



Decoupling limits of sgoldstino modes in global and local supersymmetry



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ARTICLE INFO

Article history:

Received 10 February 2013

Received in revised form 19 April 2013

Accepted 3 June 2013

Available online 14 June 2013

Editor: L. Alvarez-Gaumé

ABSTRACT

We study the decoupling limit of a superheavy sgoldstino field in spontaneously broken $\mathcal{N} = 1$ supergravity. Our approach is based on Kähler superspace, which, among others, allows direct formulation of $\mathcal{N} = 1$ supergravity in the Einstein frame and correct identifications of mass parameters. Allowing for a non-renormalizable Kähler potential in the hidden sector, the decoupling limit of a superheavy sgoldstino is identified with an infinite negative Kähler curvature. Constraints that lead to non-linear realizations of supersymmetry emerge as consequence of the equations of motion of the goldstino superfield when considering the decoupling limit. Finally, by employing superspace Bianchi identities, we identify the real chiral superfield, which will be the superconformal symmetry breaking chiral superfield that enters the conservation of the Ferrara–Zumino multiplet in the field theory limit of $\mathcal{N} = 1$ supergravity.

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1. Introduction

Supersymmetry is one of the most appealing candidates for new physics. It has not been observed so far and thus, it should be broken at some high energy scale if it is realized at all. However, supersymmetry breaking is not an easy task. In the MSSM for example, supersymmetry breaking is employed by introducing soft breaking terms. These terms are *ad hoc* masses for the superpartners of the SM particles, which nevertheless do not spoil the UV properties of the theory. In fact the MSSM includes all these soft breaking terms and one has to fit them into the observations. From a more theoretical point of view, the origin of these soft terms should be explored. The common lore is that supersymmetry should be broken in a sector of the theory, not directly connected to the SM particles, the hidden sector. For a review on soft terms, and other supersymmetry breaking mediation scenarios we refer to [1–3].

Whatever the nature of the mediation, the hidden sector should be studied on its own right. If it is a chiral multiplet that breaks supersymmetry, its highest component F will acquire a non-vanishing vev. There is a number of different scenarios for the origin of the supersymmetry breaking [1,3]. Let us note that higher derivative operators [4–7] may play an important role in hidden sector supersymmetry breaking. One of the most efficient methods

for studying the phenomenology of the hidden sector is through the dynamics of the goldstino [8–23]. The latter is the fermionic component of the superfield that breaks supersymmetry. If the supersymmetry breaking scale is low, goldstino dynamics become increasingly important for low energy phenomenology [24–35]. In fact, if the SUSY breaking scale \sqrt{f} is low with respect to Planck mass M_P ($\sqrt{f} \ll M_P$) as in gauge mediation, transverse gravitino couplings are of order M_P^{-1} and therefore are suppressed with respect to longitudinal gravitino couplings, which are of order $f^{-1/2}$. In this case, in the gravity decoupling limit, only the longitudinal gravitino component, i.e., the goldstino survives. Moreover, the highest component of the superfield to which the goldstino belongs, acquires a vev and breaks spontaneous the supersymmetry giving also mass to the sgoldstino (goldstino's superpartner). Therefore, at low energies, supersymmetry is spontaneously broken and after decoupling the sgoldstino (by making the latter superheavy) we are left with only the goldstino in the spectrum and a non-linear realized SUSY. In the case of local supersymmetry, non-linear realizations are less studied in the supergravity context [11, 42, 43].

Recently new methods have been proposed in order to study goldstino couplings, and MSSM extensions that incorporate them have been constructed [32–39]. All this framework is based on the idea of constrained superfields [10, 11, 14] that introduce a non-linear supersymmetry representation for the goldstino when its massive scalar superpartner is heavy and can be integrated out. Moreover, when one studies physics much lower than the MSSM soft masses scale, non-linear supersymmetry is realized on the SM

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particles as well, via the appropriate constraints. The constraint that enforces a non-linear supersymmetry realization for the goldstino reads

$$\Phi_{NL}^2 = 0. \quad (1.1)$$

In addition, it has been proven in [32] that in fact Φ_{NL} is proportional in the IR limit to the chiral superfield X that sources the violation of the conservation of the Ferrara–Zumino supercurrent $J_{\alpha\dot{\alpha}}$ [40,41]

$$\bar{D}^{\dot{\alpha}} J_{\alpha\dot{\alpha}} = D_{\alpha} X. \quad (1.2)$$

We extend this to the case of $\mathcal{N} = 1$ supergravity by identifying the superfield, which turns out to be the chiral superfield X of (1.2) in the gravity decoupling limit. Here, the conservation of the Ferrara–Zumino multiplet $J_{\alpha\dot{\alpha}}$ in (1.2) is replaced by the consistency conditions of the Bianchi identities [45]

$$X_{\alpha} = D_{\alpha} \mathcal{R} - \bar{D}^{\dot{\alpha}} G_{\alpha\dot{\alpha}} \quad (1.3)$$

where $G_{\alpha\dot{\alpha}}$ and \mathcal{R} are the usual supergravity superfields and $X_{\alpha} = -\frac{1}{8}(\bar{D}^2 - 8\mathcal{R})D_{\alpha} K$ is the matter sector contribution.

2. Supergravity in Einstein frame

In the standard $\mathcal{N} = 1$ superspace formulation of supergravity, one is forced to perform a Weyl rescaling to the action in order to write the theory in the Einstein frame. Here, we should write the superspace action directly in the Einstein frame since we want to correctly identify the masses to be sent to infinity. This will provide the superfield equations of motion in the correct frame as well. The appropriate framework for this is the Kähler superspace formalism which we will briefly present below. For a detailed description, one may consult for example [45–47]. An alternative method would be a super-Weyl invariant reformulation of the old minimal formulation for $\mathcal{N} = 1$ SUGRA [48].

In the conventional superspace approach to supergravity, the Lagrangian describing gravity coupled to matter would be (ignoring superpotential for the moment)

$$\mathcal{L}_F = \int d^2\theta \, 2\mathcal{E} \left\{ \frac{3}{8}(\bar{D}\bar{D} - 8\mathcal{R})e^{-\frac{1}{3}K(\Phi, \bar{\Phi})} \right\} + \text{h.c.} \quad (2.1)$$

where $2\mathcal{E}$ is the superspace chiral density and the new Θ variables span only the chiral superspace. An equivalent way to write the action (2.1) is

$$\mathcal{L}_D = -3 \int d^4\theta \, E e^{-\frac{1}{3}K(\Phi, \bar{\Phi})}, \quad (2.2)$$

where now E is the full superspace density and θ are to be integrated over the full superspace. Both actions (2.1), (2.2) can equivalently be used in order to build invariant theories in superspace. Note that \mathcal{E} and E , both have the vierbein determinant in their lowest component. As usual \mathcal{R} represents the supergravity chiral superfield which contains the Ricci scalar in its highest component. Direct calculation of (2.2) in component form shows that the theory is actually expressed in an unconventional Jordan frame. Of course a Weyl rescaling may be performed in order to bring the theory in the standard Einstein frame. Nevertheless, it is possible to perform this rescaling at the superspace level by considering

$$\begin{aligned} E_M^{\prime a} &= e^{-\frac{1}{6}K(\Phi, \bar{\Phi})} E_M^a, \\ E_M^{\prime \alpha} &= e^{-\frac{1}{12}K(\Phi, \bar{\Phi})} \left[E_M^{\alpha} - \frac{i}{12} E_M^b (\epsilon\sigma_b)^{\alpha}_{\dot{\alpha}} \bar{D}^{\dot{\alpha}} K(\Phi, \bar{\Phi}) \right], \\ E_{M\dot{\alpha}}^{\prime} &= e^{-\frac{1}{12}K(\Phi, \bar{\Phi})} \left[E_{M\dot{\alpha}} - \frac{i}{12} E_M^b (\epsilon\bar{\sigma}_b)_{\dot{\alpha}}^{\alpha} D_{\alpha} K(\Phi, \bar{\Phi}) \right] \end{aligned}$$

where E_A^M is the superspace frame, containing the gravitino and the vierbein in the appropriate lowest components. This redefinition will change the structure of the whole superspace including the Bianchi identity solutions and the superspace derivatives. Most importantly, the superspace geometry will receive contributions at the same time from the matter and supergravity fields in a unified way. The Lagrangian (2.2) now becomes in the new superspace frame (erasing the primes for convenience)

$$\mathcal{L}_{D \text{ new}} = -3 \int d^4\theta \, E. \quad (2.3)$$

This form now contains the properly normalized supergravity action coupled to matter. The interested reader should consult an extensive review on the subject [45]. Since we also wish to include a superpotential, the appropriate contribution will be given by adding to (2.3) the appropriately rescaled superpotential W so that the full Lagrangian will be given by

$$\mathcal{L}_{\text{superpotential}} = -3 \int d^4\theta \, E + \left\{ \int d^4\theta \, \frac{E}{2\mathcal{R}} e^{K/2} W + \text{h.c.} \right\}. \quad (2.4)$$

In this new framework, Kähler transformations, generated by holomorphic functions F , are expressed as field dependent transformations gauged by a composite $U_K(1)$ vector B_A . The respective charge now is referred to as “chiral weight” and a superfield Φ of chiral weight $w(\Phi)$ transforms as

$$\Phi \rightarrow \Phi e^{-\frac{i}{2}w(\Phi) \text{Im } \mathcal{F}}. \quad (2.5)$$

Gauge covariant superspace derivatives are defined as

$$\mathcal{D}_A \Phi = E_A^M \partial_M \Phi + w(\Phi) B_A \Phi \quad (2.6)$$

where the composite connection superfields are

$$\begin{aligned} B_{\alpha} &= \frac{1}{4} D_{\alpha} K, & \bar{B}^{\dot{\alpha}} &= -\frac{1}{4} \bar{D}^{\dot{\alpha}} K, \\ B_a &= \frac{1}{4} (\partial_i K) D_a \Phi^i - \frac{1}{4} (\partial_{\bar{j}} K) D_a \bar{\Phi}^{\bar{j}} + \frac{3i}{2} \mathcal{G}_a \\ &\quad + \frac{i}{8} g_{i\bar{j}} \bar{\sigma}^{\dot{\alpha}\alpha} (D_{\alpha} \Phi^i) \bar{D}_{\dot{\alpha}} \bar{\Phi}^{\bar{j}}. \end{aligned}$$

All component fields are understood to be defined appropriately via projection as usual but now with the use of these Kähler superspace derivatives. It turns out that the invariant Lagrangian containing both (2.3) and (2.4) depends only on the generalized Kähler potential

$$e^G = e^{K(\Phi, \bar{\Phi})} W(\Phi) \bar{W}(\bar{\Phi}). \quad (2.7)$$

By taking into account the chiral weights of the gravity sector and performing a Kähler transformation with parameter $\mathcal{F} = \ln W$, we find that the final expression for the most general coupling of matter to supergravity is

$$\mathcal{L} = \int d^4\theta \, E \left[-3 + \frac{1}{2\mathcal{R}} e^{\frac{G}{2}} + \frac{1}{2\mathcal{R}} e^{\frac{G}{2}} \right]. \quad (2.8)$$

It should be stressed that this form of the action is completely equivalent to the standard $\mathcal{N} = 1$ superspace formulation (2.1) to which is related by appropriate redefinitions of the superspace frames.

3. Sgoldstino decoupling

We are interested in those classes of models where the sgoldstino may become superheavy and decouples from the spectrum. In this case, it plays no role in the low energy effective theory,

and its dynamics can be integrated out by its equations of motion. Essentially, in order to be able to decouple consistently the sgoldstino degrees of freedom, one has to

1. consider the sgoldstino mass as the heavier scale in the problem, and
2. find consistent solutions for the equations of motion in that limit.

This is equivalent to taking the limit of infinitely heavy sgoldstino and integrate its equations of motion, if possible, in this limit. This work has been done in component form earlier [14] and extended recently [36,37]. We will implement the above procedure in superspace, where as we will see it is quite straightforward.

To study sgoldstino decoupling in supergravity, it is helpful to consider the corresponding decoupling in global supersymmetry.

3.1. Sgoldstino decoupling in global supersymmetry

The most general single chiral globally supersymmetric superfield Lagrangian is given by

$$\mathcal{L} = \int d^4\theta K(\Phi, \bar{\Phi}) + \left\{ \int d^2\theta W(\Phi) + \text{h.c.} \right\} \quad (3.1)$$

where, $K(\Phi, \bar{\Phi})$ is the Kähler potential, a hermitian function of the chiral superfield, and $W(\Phi)$ is the superpotential, a holomorphic function of the chiral superfield. From the above action, the superspace equations of motion

$$-\frac{1}{4}\bar{D}\bar{D}K_\Phi + W_\Phi = 0, \quad (3.2)$$

with $K_\Phi = \partial_\Phi K$, $W_\Phi = \partial_\Phi W$ easily follow. For a general, non-renormalizable supersymmetric model where supersymmetry is spontaneously broken, the supertrace mass formula reads [43]

$$\text{Str } M^2 = \sum_J (-1)^{2J} (2J+1) M_J^2 = -2R_{A\bar{A}} f \bar{f} \quad (3.3)$$

where $f = \langle F \rangle$ and $R_{A\bar{A}}$ ($= g^{A\bar{A}} R_{A\bar{A}A\bar{A}}$) is the Ricci tensor of the scalar Kähler manifold evaluated at the vacuum expectation values of the scalars. Eq. (3.3) describes the mass splitting between the components of the supermultiplet. In the case of a single chiral superfield we are discussing, since the goldstino is always massless, the supertrace of the goldstino multiplet is just the square of the sgoldstino mass

$$M_{\text{sg}}^2 = -R_{A\bar{A}} f \bar{f} \quad (3.4)$$

We see that necessarily the scalar manifold should be a space of negative curvature in order to have non-tachyonic scalar excitations. In addition, the limit of the infinitely heavy sgoldstino

$$2M_{\text{sg}}^2 = \text{Str } M^2 \rightarrow \infty \quad \text{or} \quad R_{A\bar{A}A\bar{A}} \rightarrow -\infty. \quad (3.5)$$

Since

$$R_{A\bar{A}A\bar{A}} = \partial_{\bar{A}} \partial_A \partial_{\bar{A}} \partial_A K - \partial_{\bar{A}} \partial_A \partial_{\bar{A}} K \partial_A \partial_A K, \quad (3.6)$$

in normal coordinates for the Kähler space in which $g_{A\bar{A}} = \delta_{A\bar{A}}$ and $\partial_i \partial_j \partial_k K = 0$ (for any $i, j = A, \bar{A}$), we have that the infinitely heavy sgoldstino is obtained in the limit

$$-\partial_{\bar{A}} \partial_A \partial_{\bar{A}} \partial_A K \rightarrow \infty. \quad (3.7)$$

By assuming that the vacuum expectation value of $A = \Phi|$ vanishes,¹ the general form of the Kähler potential

$$K(\Phi, \bar{\Phi}) = \sum_{mn} c_{mn} \Phi^m \bar{\Phi}^n \quad (3.8)$$

will have the following expansion in normal coordinates

$$K(\Phi, \bar{\Phi}) = \Phi \bar{\Phi} + c_{22} \bar{\Phi}^2 \Phi^2 + \dots \quad (3.9)$$

It is easy to see that in fact

$$c_{22} = \frac{1}{4} R_{A\bar{A}A\bar{A}} = \frac{1}{4} R_{A\bar{A}} \quad (3.10)$$

in normal coordinates. By using then (3.3), (3.5), we get that the Kähler potential may be expressed in terms of the sgoldstino mass as

$$K(\Phi, \bar{\Phi}) = \Phi \bar{\Phi} - \frac{M_{\text{sg}}^2}{4|f|^2} \bar{\Phi}^2 \Phi^2 + \dots \quad (3.11)$$

where the dots stands for M_{sg} -independent terms and $f = \langle F \rangle$ is the vev of the auxiliary field in the chiral multiplet. From the superspace equations of motion (3.2), one can easily isolate the contribution proportional to M_{sg}^2 . Indeed, (3.2) is written as

$$\frac{M_{\text{sg}}^2}{4|f|^2} \Phi \bar{D}\bar{D}\bar{\Phi}^2 + (M_{\text{sg}}\text{-independent terms}) = 0. \quad (3.12)$$

Therefore, in the $M_{\text{sg}} \rightarrow \infty$ limit, the M_{sg} -dependent part of the field equations is turned into the superspace constraint

$$\Phi \bar{D}\bar{D}\bar{\Phi}^2 = 0. \quad (3.13)$$

To explicitly solve (3.13), we note that it leads to three component equations

$$\Phi \bar{D}\bar{D}\bar{\Phi}^2| = 0, \quad D_\alpha(\Phi \bar{D}\bar{D}\bar{\Phi}^2)| = 0, \quad DD(\Phi \bar{D}\bar{D}\bar{\Phi}^2)| = 0. \quad (3.14)$$

The non-trivial solution to the above equations is [10,32]

$$\Phi_{NL} = \frac{\chi\bar{\chi}}{2F} + \sqrt{2}\theta\chi + \theta^2 F \quad (3.15)$$

which can be easily checked that it satisfies

$$\Phi_{NL}^2 = 0. \quad (3.16)$$

As a result, the sgoldstino can be safely decoupled in the $M_{\text{sg}} \rightarrow \infty$ limit as long as Φ satisfies (3.13), or equivalently (3.16).

3.2. Sgoldstino decoupling in supergravity

As in the case of global supersymmetry, we are interested in the equations of motion and the mass supertrace. The superfield equations of motion as follow from the action (2.8) are [46]

$$\mathcal{R} = \frac{1}{2} e^{\frac{G}{2}}, \quad (3.17)$$

$$\mathcal{G}_\alpha + \frac{1}{8} G_{\phi\bar{\phi}} \bar{\sigma}_\alpha^{\dot{\alpha}\alpha} \mathcal{D}_\alpha \Phi \bar{\mathcal{D}}_{\dot{\alpha}} \bar{\Phi} = 0, \quad (3.18)$$

$$(\bar{\mathcal{D}}\bar{\mathcal{D}} - 8\mathcal{R})G_\phi = 0. \quad (3.19)$$

On the other hand, for a general supergravity model with only one chiral multiplet the supertrace is given by [44]

$$\text{Str } M^2 = -2R_{A\bar{A}} f \bar{f}, \quad (3.20)$$

which means that in the limit of infinite negative Kähler curvature the sgoldstino will become superheavy and can consistently be integrated out. Indeed, (3.20) is explicitly written as

¹ if not we may shift appropriately A so that $\langle A \rangle = 0$

$$M_{\text{sg}}^2 = 2m_{3/2}^2 - R_{A\bar{A}}\bar{f}\bar{f}. \quad (3.21)$$

Therefore, for finite gravitino mass $m_{3/2}$, the infinite curvature limit

$$R_{A\bar{A}\bar{A}\bar{A}} \rightarrow -\infty \quad (3.22)$$

is equivalent to superheavy sgoldstinos. Again, in normal coordinates

$$R_{A\bar{A}\bar{A}\bar{A}} = \partial_{\bar{A}}\partial_A\partial_{\bar{A}}\partial_A K = \partial_{\bar{A}}\partial_A\partial_{\bar{A}}\partial_A G \quad (3.23)$$

and therefore with

$$G \supset \frac{2m_{3/2}^2 - M_{\text{sg}}^2}{4|f|^2} \phi^2 \bar{\phi}^2 + \dots \quad (3.24)$$

the decoupling limit we are after is again $M_{\text{sg}}^2 \rightarrow \infty$. Taking into account that the Kähler curvature $M_{\text{sg}}^2/4|f|^2$ will dominate the equations of motion and following the same reasoning as in the global supersymmetric case, we get from (3.19)

$$\Phi(\bar{D}\bar{D} - 8\mathcal{R})\bar{\phi}^2 = 0. \quad (3.25)$$

This constraint is the curved superspace analogue of (3.13). In order to solve it, we take into account that $\Phi(\bar{D}\bar{D} - 8\mathcal{R})\bar{\phi}^2$ is a chiral superfield, and we will once again start from its lowest component, namely

$$\Phi(\bar{D}\bar{D} - 8\mathcal{R})\bar{\phi}^2| = 0. \quad (3.26)$$

This is written, for

$$\Phi = A + \sqrt{2}\Theta\chi + \Theta\Theta F, \quad \mathcal{R}| = -\frac{1}{6}M \quad (3.27)$$

as

$$AM\bar{A}^2 - 24A\bar{A}\bar{F} + 12A\bar{\chi}\bar{\chi} = 0. \quad (3.28)$$

This equation has three solutions

$$A_0 = 0, \quad A_1 = \frac{\chi\chi}{2F}, \quad A_2 = \frac{24F}{M} - \frac{\chi\chi}{2F}. \quad (3.29)$$

The first solution A_0 is the trivial and we will not consider it. The second solution A_1 is the $\phi^2 = 0$ we already encounter in the global SUSY case. The third solution A_2 corresponds to $\phi^2 \neq 0$ and can only be realized as long as the auxiliary field of supergravity M is non vanishing ($M \neq 0$). However, from Eq. (3.17) we get

$$\mathcal{R} = \frac{1}{2}e^{\frac{G}{2}} = \frac{1}{2}e^{-\frac{M_{\text{sg}}^2}{8|f|^2}\phi^2\bar{\phi}^2+\dots}, \quad (3.30)$$

where only the dominant term was explicitly written in the exponent in the right hand side. Now, in the $M_{\text{sg}}^2 \rightarrow \infty$ limit, the right hand side goes to zero exponentially fast so that for $\phi^2 \neq 0$

$$\mathcal{R} = 0 \quad \text{for } M_{\text{sg}}^2 \rightarrow \infty. \quad (3.31)$$

Therefore also $M = -6\mathcal{R}| = 0$ and the third solution (A_2) cannot consistently be realized. As a result, the only solution to the constraint (3.25) is the $A_1 = \frac{\chi\chi}{2F}$, or in other words the familiar

$$\phi^2 = 0. \quad (3.32)$$

This constraint leads to

$$e^{\frac{M_{\text{sg}}^2}{8|f|^2}\phi^2\bar{\phi}^2}\Big|_{\phi^2=0} = 1 \quad (3.33)$$

and thus, the divergent part of (3.17) completely decouples! Moreover, $\phi^2 = 0$ also satisfies

$$\mathcal{D}_\alpha\Phi\bar{\mathcal{D}}_{\dot{\alpha}}\bar{\phi}^2 = 0 \quad (3.34)$$

which is the field equation (3.18) in the $M_{\text{sg}}^2 \rightarrow \infty$ limit. As a result, we have again arrived to the constraint (3.32) as the only viable and consistent condition for the decoupling of the sgoldstino.

3.3. Supercurrent and sgoldstino decoupling

In order to discuss the relation of supersymmetry breaking to conservation laws, let us explore the decoupling limit of the supergravity sector. The supergravity equations of motion (3.17) and (3.18) in superspace, after restoring dimensions with compensating powers of M_P and returning to the Kähler frame where everything is expressed in terms of K and W , are written as

$$\mathcal{R} = \frac{1}{M_P^2} \frac{1}{2} W e^{\frac{K}{2M_P^2}}, \quad (3.35)$$

$$\mathcal{G}_a + \frac{1}{M_P^2} \frac{1}{8} g_{i\bar{j}} \bar{\sigma}_a^{\dot{\alpha}\alpha} \mathcal{D}_\alpha \Phi^i \bar{\mathcal{D}}_{\dot{\alpha}} \bar{\phi}^{\bar{j}} = 0. \quad (3.36)$$

Gravity decouples in the limit $M_P \rightarrow \infty$, and from (3.35) and (3.36) we have

$$\mathcal{R} \rightarrow 0, \quad \mathcal{G}_a \rightarrow 0. \quad (3.37)$$

We note that this is the limit even when $W/M_P = \text{finite}$, which is another possible limit [42] for gauge mediated SUSY breaking scenarios. The fact that these supergravity superfields should vanish can be also understood from the algebra of supergravity when compared to supersymmetry. For example, the global commutation relation (for $w(\Phi^i) = 0$)

$$[\bar{D}_{\dot{\alpha}}, D_a]\Phi^i = 0, \quad (3.38)$$

in supergravity becomes

$$[\bar{D}_{\dot{\alpha}}, D_a]\Phi^i = -i\mathcal{R}\sigma_{a\dot{\alpha}}^\alpha \mathcal{D}^\alpha \Phi^i \quad (3.39)$$

thus in order to recover the global supersymmetry algebra the superfield \mathcal{R} should vanish.

Let us now derive the analog of the conservation equation of the Ferrara-Zumino multiplet (1.2) in curved superspace. By using the consistency conditions of the Bianchi identities [45]

$$X_\alpha = M_P^2 \mathcal{D}_\alpha \mathcal{R} - M_P^2 \bar{\mathcal{D}}^{\dot{\alpha}} G_{\alpha\dot{\alpha}} \quad (3.40)$$

with

$$X_\alpha = -\frac{1}{8}(\bar{D}^2 - 8\mathcal{R})\mathcal{D}_\alpha K \quad (3.41)$$

and the equations of motion, we find

$$\bar{\mathcal{D}}^{\dot{\alpha}} \mathcal{J}_{\alpha\dot{\alpha}} = \mathcal{D}_\alpha \mathcal{X} - \frac{16}{3} \mathcal{R} \mathcal{D}_\alpha K + \frac{2}{3} G_{\alpha\dot{\alpha}} \bar{\mathcal{D}}^{\dot{\alpha}} K \quad (3.42)$$

with

$$\begin{aligned} \mathcal{J}_{\alpha\dot{\alpha}} &= 2g_{i\bar{j}} \mathcal{D}_\alpha \Phi^i \bar{\mathcal{D}}_{\dot{\alpha}} \bar{\phi}^{\bar{j}} - \frac{2}{3} [\mathcal{D}_\alpha, \bar{\mathcal{D}}_{\dot{\alpha}}] K, \\ \mathcal{X} &= 4W e^{\frac{K}{2M_P^2}} - \frac{1}{3} \bar{\mathcal{D}} \bar{\mathcal{D}} K. \end{aligned} \quad (3.43)$$

The extra terms compared to (1.2) arise due to commutation relations like (3.39), and should vanish when supergravity is decoupled.

Now we take the decoupling limit of supergravity ($M_P \rightarrow \infty$) with ($\mathcal{R} \rightarrow 0$, $\mathcal{G}_a \rightarrow 0$) and find exactly the same formula as the global case. As a final comment let us note that now, after the decoupling of supergravity, the superfield X is

$$\mathcal{X} \rightarrow X = 4W - \frac{1}{3} \bar{D} \bar{D} K. \quad (3.44)$$

4. Conclusions

In this work we explored the decoupling limit of sgoldstinos in spontaneously broken SUSY theories. This decoupling was implemented by considering large mass values for the sgoldstino (in fact the infinite mass limit). We used superspace techniques as they allowed for a unified treatment of the spontaneous breaking of SUSY both in local and global supersymmetric cases. The motivation of this study was twofold: first to check if the constraint superfield formalism employed in the global supersymmetry still works in supergravity as well and second, to correctly identify in supergravity the chiral superfield that enters in the conservation of the Ferrara–Zumino multiplet and which accommodates the goldstino in global supersymmetry.

The way to approach these targets was to reformulate the goldstino dynamics in global supersymmetry but now in a language appropriate for supergravity. First we have identified the sgoldstino mass in both cases, and found the decoupling limit (supermassive sgoldstino) to be the limit of infinite negative Kähler curvature. Then we impose this limit to the superfield equations of motion and in the case of supersymmetry we found the constraint ($\Phi \bar{D}^2 \bar{\Phi}^2 = 0$) which is solved by $\Phi^2 = 0$ as expected. In the case of supergravity, the super-covariant form of the more general constraint emerges, but again with the same single consistent solution. Thus, the superspace constraint $\Phi^2 = 0$ for the goldstino, when the sgoldstino is supermassive, holds for supergravity as well. However, we should mention a potential problem here. Namely, the expansion of the Kähler potential in (3.11) is written in powers of M_{sg}/f , from where it follows that actually $M_{sg} \sim f/\Lambda$ where Λ is the effective cutoff of the theory. The infinite sgoldstino mass seems therefore to be in conflict with the removal of the cutoff ($\Lambda \rightarrow \infty$), which is needed to identify the goldstino superfield with the infrared limit of the superconformal symmetry breaking superfield that enters the Ferrara–Zumino current conservation. This issue is further complicated by the presence of extra light fields. The problem has been pointed out in [38] where conditions for the effective expansion of the supersymmetric Lagrangian in terms of the inverse cutoff to not be in conflict with a small sgoldstino mass $\sim f/\Lambda$ were given. Note that we have not faced this problem, as we have taken the formal infinite large sgoldstino mass limit.

Acknowledgements

We would like to thank Ulf Lindström for discussions and comments. This research was implemented under the “ARISTEIA” Action of the “Operational Programme Education and Lifelong Learning” and is co-funded by the European Social Fund (ESF) and National Resources.

References

- [1] D.Z. Freedman, A. Van Proeyen, *Supergravity*, Cambridge Univ. Press, Cambridge, UK, 2012, 607 p.
- [2] S. Weinberg, *The Quantum Theory of Fields. Vol. 3: Supersymmetry*, Cambridge Univ. Press, Cambridge, UK, 2000, 419 p.

- [3] S.P. Martin, A supersymmetry primer, in: G.L. Kane (Ed.), *Perspectives on Supersymmetry II*, 2010, pp. 1–153, arXiv:hep-ph/9709356.
- [4] M. Koehn, J.-L. Lehners, B.A. Ovrut, Higher-derivative chiral superfield actions coupled to $N = 1$ supergravity, *Phys. Rev. D* 86 (2012) 085019, arXiv:1207.3798 [hep-th].
- [5] M. Koehn, J.-L. Lehners, B.A. Ovrut, The ghost condensate in $N = 1$ supergravity, arXiv:1212.2185 [hep-th].
- [6] F. Farakos, A. Kehagias, Emerging potentials in higher-derivative gauged chiral models coupled to $N = 1$ supergravity, arXiv:1207.4767 [hep-th].
- [7] S. Cecotti, S. Ferrara, L. Girardello, Structure of the scalar potential in general $N = 1$ higher derivative supergravity in four-dimensions, *Phys. Lett. B* 187 (1987) 321.
- [8] D.V. Volkov, V.P. Akulov, Is the neutrino a Goldstone particle?, *Phys. Lett. B* 46 (1973) 109.
- [9] E.A. Ivanov, A.A. Kapustnikov, General relationship between linear and nonlinear realizations of supersymmetry, *J. Phys. A* 11 (1978) 2375.
- [10] M. Rocek, Linearizing the Volkov–Akulov model, *Phys. Rev. Lett.* 41 (1978) 451.
- [11] U. Lindstrom, M. Rocek, Constrained local superfields, *Phys. Rev. D* 19 (1979) 2300.
- [12] T. Uematsu, C.K. Zachos, Structure of phenomenological Lagrangians for broken supersymmetry, *Nucl. Phys. B* 201 (1982) 250.
- [13] Q. Ho-Kim, A. Smailagic, Constrained superfields, *Phys. Rev. D* 29 (1984) 2427.
- [14] R. Casalbuoni, S. De Curtis, D. Dominicis, F. Feruglio, R. Gatto, Nonlinear realization of supersymmetry algebra from supersymmetric constraint, *Phys. Lett. B* 220 (1989) 569.
- [15] M.A. Luty, E. Ponton, Effective Lagrangians and light gravitino phenomenology, *Phys. Rev. D* 57 (1998) 4167, arXiv:hep-ph/9706268.
- [16] T. Lee, G.H. Wu, Nonlinear supersymmetric effective Lagrangian and goldstino interactions at high energies, *Mod. Phys. Lett. A* 13 (1998) 2999, arXiv:hep-ph/9811458.
- [17] A. Brignole, J.A. Casas, J.R. Espinosa, I. Navarro, Low-scale supersymmetry breaking: Effective description, electroweak breaking and phenomenology, *Nucl. Phys. B* 666 (2003) 105, arXiv:hep-ph/0301121.
- [18] I. Antoniadis, M. Tuckmantel, Non-linear supersymmetry and intersecting D-branes, *Nucl. Phys. B* 697 (2004) 3, arXiv:hep-th/0406010.
- [19] I. Antoniadis, M. Tuckmantel, F. Zwirner, Phenomenology of a leptonic goldstino and invisible Higgs boson decays, *Nucl. Phys. B* 707 (2005) 215, arXiv:hep-ph/0410165.
- [20] J.A. Bagger, A.F. Falk, Decoupling and destabilizing in spontaneously broken supersymmetry, *Phys. Rev. D* 76 (2007) 105026, arXiv:0708.3364 [hep-ph].
- [21] Z. Komargodski, N. Seiberg, Comments on the Fayet–Iliopoulos term in field theory and supergravity, *JHEP* 0906 (2009) 007, arXiv:0904.1159 [hep-th].
- [22] S.M. Kuzenko, S.J. Tyler, Complex linear superfield as a model for Goldstino, *JHEP* 1104 (2011) 057, arXiv:1102.3042 [hep-th].
- [23] S.M. Kuzenko, S.J. Tyler, On the Goldstino actions and their symmetries, *JHEP* 1105 (2011) 055, arXiv:1102.3043 [hep-th].
- [24] C. Petersson, A. Romagnoni, The MSSM Higgs sector with a dynamical goldstino supermultiplet, *JHEP* 1202 (2012) 142, arXiv:1111.3368 [hep-ph].
- [25] B. Bellazzini, C. Petersson, R. Torre, Photophilic Higgs from sgoldstino mixing, arXiv:1207.0803 [hep-ph].
- [26] H.-C. Cheng, W.-C. Huang, I. Low, A. Menon, Goldstini as the decaying dark matter, *JHEP* 1103 (2011) 019, arXiv:1012.5300 [hep-ph].
- [27] T.E. Clark, S.T. Love, T. ter Veldhuis, Supersymmetric multi-Higgs doublet model with non-linear electroweak symmetry breaking, *Phys. Rev. D* 85 (2012) 015014, arXiv:1107.3116 [hep-ph].
- [28] R. Argurio, Z. Komargodski, A. Mariotti, Pseudo-goldstini in field theory, *Phys. Rev. Lett.* 107 (2011) 061601, arXiv:1102.2386 [hep-th].
- [29] L. Alvarez-Gaumé, C. Gomez, R. Jimenez, Phenomenology of the minimal inflation scenario: inflationary trajectories and particle production, *JCAP* 1203 (2012) 017, arXiv:1110.3984 [astro-ph.CO].
- [30] L. Alvarez-Gaumé, C. Gomez, R. Jimenez, A minimal inflation scenario, *JCAP* 1103 (2011) 027, arXiv:1101.4948 [hep-th].
- [31] L. Alvarez-Gaumé, C. Gomez, R. Jimenez, Minimal inflation, *Phys. Lett. B* 690 (2010) 68, arXiv:1001.0010 [hep-th].
- [32] Z. Komargodski, N. Seiberg, From linear SUSY to constrained superfields, *JHEP* 0909 (2009) 066, arXiv:0907.2441 [hep-th].
- [33] I. Antoniadis, E. Dudas, D.M. Ghilencea, P. Tziveloglou, Non-linear MSSM, *Nucl. Phys. B* 841 (2010) 157, arXiv:1006.1662 [hep-ph].
- [34] I. Antoniadis, E. Dudas, D.M. Ghilencea, P. Tziveloglou, Nonlinear supersymmetry and goldstino couplings to the MSSM, *Theor. Math. Phys.* 170 (2012) 26, *Theor. Mat. Fiz.* 170 (2012) 34.
- [35] F. Farakos, A. Kehagias, Non-linear single Higgs MSSM, arXiv:1210.4941 [hep-ph].
- [36] E. Dudas, G. von Gersdorff, D.M. Ghilencea, S. Lavignac, J. Parmentier, On non-universal Goldstino couplings to matter, *Nucl. Phys. B* 855 (2012) 570, arXiv:1106.5792 [hep-th].
- [37] I. Antoniadis, E. Dudas, D.M. Ghilencea, Goldstino and sgoldstino in microscopic models and the constrained superfields formalism, *Nucl. Phys. B* 857 (2012) 65, arXiv:1110.5939 [hep-th].

- [38] I. Antoniadis, D.M. Ghilencea, Low-scale SUSY breaking and the (s)goldstino physics, arXiv:1210.8336 [hep-th].
- [39] E. Dudas, C. Petersson, P. Tziveloglou, Low scale supersymmetry breaking and its LHC signatures, arXiv:1211.5609 [hep-ph].
- [40] S. Ferrara, B. Zumino, Transformation properties of the supercurrent, Nucl. Phys. B 87 (1975) 207.
- [41] T.E. Clark, S.T. Love, The supercurrent in supersymmetric field theories, Int. J. Mod. Phys. A 11 (1996) 2807, arXiv:hep-th/9506145.
- [42] S.P. de Alwis, A local evaluation of global issues in SUSY breaking, arXiv:1211.3913 [hep-th].
- [43] S.J. Gates, M.T. Grisaru, M. Rocek, W. Siegel, Superspace or one thousand and one lessons in supersymmetry, Front. Phys. 58 (1983) 1, arXiv:hep-th/0108200.
- [44] J. Wess, J. Bagger, Supersymmetry and Supergravity, Princeton Univ. Press, Princeton, USA, 1992, p. 259.
- [45] P. Binetrui, G. Girardi, R. Grimm, Supergravity couplings: A geometric formulation, Phys. Rep. 343 (2001) 255, arXiv:hep-th/0005225.
- [46] P. Binetrui, G. Girardi, R. Grimm, M. Muller, Kahler transformations and the coupling of matter and Yang–Mills fields to supergravity, Phys. Lett. B 189 (1987) 83.
- [47] R. Grimm, Supergravity, Kahler superspace and gauged supersymmetric sigma models, Phys. Lett. B 242 (1990) 64.
- [48] S.M. Kuzenko, S.A. McCarthy, On the component structure of $N = 1$ supersymmetric nonlinear electrodynamics, JHEP 0505 (2005) 012, arXiv:hep-th/0501172.