Deformed N=1 supersymmetry

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Abstract

We consider a deformation of N=1 four dimensional Minkowski superspace where odd coordinates $\theta^\alpha$ do not anticommute. We define supersymmetric and associative star product and show how the remaining (anti)commutation relations among the superspace coordinates are modified. In particular, the even coordinates do not commute as well. We also study chiral and vector superfields and their interactions. Surprisingly we find that ordinary undeformed N=1 supersymmetric field theories are compatible with the deformed supersymmetry considered.
1 Introduction

It is widely believed that further progress in our understanding of elementary particles and fundamental interactions is ultimately related with deeper understanding of the actual structure of space-time at short distances. The idea that usual four space-time coordinates could be supplemented by anticommuting spinorial coordinates \( \theta^\alpha, \bar{\theta}^\dot{\alpha} \),

\[
\{ \theta^\alpha, \theta^\beta \} = \{ \bar{\theta}^{\dot{\alpha}}, \bar{\theta}^{\dot{\beta}} \} = 0, \quad (1)
\]

\[
\{ \theta^\alpha, \bar{\theta}^{\dot{\alpha}} \} = 0, \quad (2)
\]

leads to a notion of (N=1) superspace. Field theories on such extension of space-time have improved ultra-violet (UV) behavior compared to the usual fields theories due to the remarkable symmetry between the fields with Bose and Fermi statistics, called supersymmetry (SUSY). Largely because of this N=1 SUSY is currently considered as a prime candidate for the fundamental particles and interactions beyond the celebrated Standard Model.

Another interesting theoretical approach, which originally was also thought to be responsible for the curing of UV divergences, is space-time noncommutativity. In the past few years noncommutative field theories (see e.g. [1] for a review) have attracted enormous interest after the realization that the noncommutativity arises naturally as a low energy limit of string theory in the background of constant antisymmetric NS-NS B-field [2]. However, noncommutative field theories show up some pathological features which prevent to apply them to particle physics (see e.g. [3] and [4] for the attempts to construct noncommutative Standard Model). The most difficult problem is the occurrence of infrared divergences [5] which perhaps indicate that pure field theoretical description is incomplete. Other difficulties are related to peculiar properties of noncommutative gauge groups and their representations which are reflected in charge quantization problem [6] and to anomalies [7] which make construction of desired chiral theories problematic.

Straightforward SUSY extension of noncommutative field theories where the odd superspace coordinates satisfy the usual (anti)commutation relations [8] is of little help in this situation. However, it has been also realized that SUSY is actually compatible with more general (anti)commutation relations [9, 10, 11]. More recently, non-trivial (anti)commutation relations for the odd coordinates has been obtained from the string theory in the constant graviphoton and/or gravitino backgrounds [12, 13, 14]. This observation initiates subsequently recent studies [14, 15, 16, 17] of field theories with deformed SUSY mainly in Euclidean space.\(^1\)

Since our main interest is possible implications of deformed SUSY in particle physics, in the present paper we study a particular deformation of N=1 four dimensional Minkowski superspace. In particular, instead of ordinary anticommutation relations (1) we assume that

\^1After the completion of this work recent paper [18] appeared which considers the general case of deformed Poisson brackets including N=1 supersymmetric case.
odd superspace coordinates do not anticommute. In subsequent sections we define supersymmetric and associative Weyl-Moyal-type star products and show how the consistency dictates modification of (anti)commutation relations of other superspace coordinates. In particular, the even coordinates do not commute as well. Then we define the relevant superfields and construct invariant actions. Somewhat surprisingly we will find that ordinary undeformed supersymmetric field theories are compatible with the deformed N=1 supersymmetry considered.

2 Deformed superspace

Consider N=1 four dimensional Minkowski superspace with a set of coordinates \( (x^\mu, \theta^\alpha, \bar{\theta}^\dot{\alpha}) \). We would like to consider deformation of anticommutation relations (1). Namely, assume

\[
\{ \hat{\theta}^\alpha, \hat{\theta}^\beta \} = C^{\alpha\beta}, \tag{3}
\]

where \( C^{\alpha\beta} = C^{\beta\alpha} \) are complex constants. Since in Minkowski space-time \( (\hat{\theta}^\alpha)^\dagger = \bar{\theta}^\dot{\alpha} \), we also have

\[
\{ \bar{\theta}^{\dot{\alpha}}, \bar{\theta}^{\dot{\beta}} \} = \bar{C}^{\dot{\alpha}\dot{\beta}}, \tag{4}
\]

\( \bar{C}^{\dot{\alpha}\dot{\beta}} = (C^{\beta\alpha})^\dagger \). Because of (3), product of functions of \( \hat{\theta}^\alpha \) should be correspondingly ordered. As in the ordinary noncommutative case, this can be done by a suitably defined Weyl-Moyal-type star product of functions of ordinary anticommuting \( \theta^\alpha \)’s. We define (see also [18]):

\[
f (\hat{\theta}) g (\hat{\theta}) \equiv f (\theta) \star g (\theta) \overset{\text{def}}{=} f (\theta) \exp \left[ -\overleftarrow{D}_\alpha \frac{C^{\alpha\beta}}{2} \overrightarrow{D}_\beta \right] g (\theta)
\]

\[
= f (\theta) \left[ 1 - \overleftarrow{D}_\alpha \frac{C^{\alpha\beta}}{2} \overrightarrow{D}_\beta - \frac{1}{16} \det \overleftarrow{C} \overleftarrow{D}^2 \overrightarrow{D}^2 \right] g (\theta), \tag{5}
\]

where \( \overleftarrow{D}_\alpha \) and \( \overrightarrow{D}_\alpha \) are left and right covariant derivatives, respectively,

\[
\overleftarrow{D}_\alpha = \overleftarrow{\partial}_\alpha + i \sigma^\mu_{\alpha\alpha} \overleftarrow{\partial}_\mu, \quad \overrightarrow{D}_\alpha = -\overrightarrow{\partial}_\alpha + i \sigma^\mu_{\alpha\alpha} \overrightarrow{\partial}_\mu \tag{6}
\]

and \( \overleftarrow{D}^2 = 2 \overleftarrow{D}_2 \overrightarrow{D}_1, \overrightarrow{D}^2 = 2 \overrightarrow{D}_1 \overleftarrow{D}_2 \). Then we have \( \{ \theta^\alpha, \theta^\beta \}_\star = C^{\alpha\beta} \). Here and below the subscript \( \star (\overline{\star}) \) means that \( \star \)-product (\( \overline{\star} \)-product) is involved. Taking Hermitian conjugate of (5) we define conjugate star product of functions of \( \bar{\theta}^{\dot{\alpha}} \):

\[
g (\bar{\theta}) f (\bar{\theta}) \equiv g (\theta) \overline{\star} f (\theta) \overset{\text{def}}{=} g (\theta) \exp \left[ -\overleftarrow{D}_\dot{\alpha} \frac{C^{\dot{\alpha}\dot{\beta}}}{2} \overrightarrow{D}_\beta \right] f (\theta)
\]

\(^2\text{We follow the conventions of Wess and Bagger [19].}\)
\[ f(\hat{\theta}) g(\hat{\theta}) = f(\theta) g(\theta). \] (9)

Obviously, the covariant derivatives satisfy the following anticommutation relations

\[ \{ \overrightarrow{D}_\alpha, \overrightarrow{D}_\beta \}_\times = 0, \quad \{ \overrightarrow{D}_\alpha, \overrightarrow{D}_\beta \}_\tau = 0, \] (10)

\[ \{ \overrightarrow{D}_\alpha, \overrightarrow{D}_\beta \}_\star = -2i\sigma^\mu_{\alpha\beta} \frac{\partial}{\partial x^\mu}. \] (11)

However,

\[ \{ \overrightarrow{D}_\alpha, \overrightarrow{D}_\beta \}_\tau = \sigma^\mu_{\alpha\beta} \partial_{x^\mu} C^{\beta\gamma} \frac{\partial}{\partial x^\gamma}, \quad \{ \overrightarrow{D}_\alpha, \overrightarrow{D}_\beta \}_\star = C^{\alpha\beta} \sigma^\mu_{\alpha\beta} \sigma^{\nu\gamma} \frac{\partial}{\partial x^\mu} \frac{\partial}{\partial x^\nu}. \] (12)

This means that \( D_\alpha(\overrightarrow{D}_\alpha) \) does not act as a derivation with respect to \( \overline{\tau}(\star) \)-product, but it does with respect to \( \star(\overline{\tau}) \)-product. Then it follows that the subalgebra of chiral (antichiral) superfields is not closed with respect to \( \star(\overline{\tau}) \)-product.

The star products (5) and (7) are associative and supersymmetric [18]. The associativity can be directly verified using (5) and (7) and it actually is the result of anticommutation relations (10). Also one can easily see that the supercharges

\[ \overrightarrow{Q}_\alpha = \overrightarrow{D}_\alpha - 2i\sigma^\mu_{\alpha\beta} \frac{\partial}{\partial x^\mu} \] (13)

\[ \overrightarrow{Q}_{\dot{\alpha}} = \overrightarrow{D}_{\dot{\alpha}} + 2i\theta^\alpha \sigma^\mu_{\alpha\beta} \frac{\partial}{\partial x^\mu}, \] (14)

act as a derivations with respect to both \( \star \) and \( \overline{\tau} \) products as a result of anticommutation relations:

\[ \{ \overrightarrow{Q}_\alpha, \overrightarrow{D}_\beta \}_\star = \{ \overrightarrow{Q}_{\dot{\alpha}}, \overrightarrow{D}_\beta \}_\star = \{ \overrightarrow{Q}_\alpha, \overrightarrow{D}_{\dot{\beta}} \}_\tau = \{ \overrightarrow{Q}_{\dot{\alpha}}, \overrightarrow{D}_{\dot{\beta}} \}_\tau = 0. \] (15)
Having defined the above star products, the remaining (anti)commutation relations among the superspace coordinates follow immediately:

\[ \star \text{- deformation : } \quad \{ \theta^\alpha, \theta^\beta \} = 0, \quad \{ \overline{\theta}^\alpha, \overline{\theta}^\beta \} = 0, \]
\[ [x^\mu, \theta^\alpha]_\star = iC^{\alpha\beta} \sigma^\mu_{\beta\gamma} \overline{\theta}^\gamma, \quad [x^\mu, \overline{\theta}^\beta] = 0, \quad (16) \]
\[ [x^\mu, x^\nu]_\star = C^{\alpha\beta}\sigma^\mu_{\alpha\gamma} \varepsilon_{\gamma\beta} \overline{\theta} \overline{\theta} \]

\[ \overline{\star} \text{- deformation : } \quad \{ \overline{\theta}^\alpha, \overline{\theta}^\beta \overline{\theta} \} = 0, \quad \{ \theta^\alpha, \theta^\beta \} \overline{\theta} = 0, \]
\[ [x^\mu, \theta^\alpha]_{\overline{\star}} = 0, \quad [x^\mu, \overline{\theta}^\beta]_{\overline{\star}} = i\overline{\theta}^\beta \sigma^\mu_{\beta\gamma} \overline{C}^{\gamma\alpha}, \quad (17) \]
\[ [x^\mu, x^\nu]_{\overline{\star}} = \theta \overline{\theta} \varepsilon_{\alpha\gamma} \sigma^\mu_{\alpha\beta} \overline{C}^{\gamma\beta}. \]

In particular one sees that even coordinates \( x^\mu \) do not commute as well.

For a product of generic superfields we can adopt either \( \star \)-product or \( \overline{\star} \)-product. Peculiar property of the above star products is that their are not Hermitian, i.e.

\[ \left( f(x, \theta, \overline{\theta}) \star g(x, \theta, \overline{\theta}) \right) \dagger = \overline{g(x, \theta, \overline{\theta})} \overline{\star} f(x, \theta, \overline{\theta}). \quad (18) \]

However, the hermiticity of Lagrangians ensures that both star deformations are physically equivalent. In what follows in the rest of the paper we adopt \( \star \)-deformation for definiteness.

### 3 Chiral superfields

The chiral (anticontiral) superfield is defined to satisfy \( \overline{D}_\alpha \Phi = 0 \) \( (\overline{D}_{\alpha} \overline{\Phi} = 0) \). In the case of undeformed supersymmetry chiral (anticontiral) superfields form a close subalgebra, i.e. any product of chiral (antichiral) superfields is a chiral (antichiral) superfield. However, if we consider \( \star \)-deformed multiplication this is no longer true. Namely, as it can be easily checked using (5), chiral superfields do not form the closed subalgebra under the \( \star \)-product, while antichiral superfields do. \textit{Vice versa}, a \( \overline{\star} \)-product of antichiral superfields is not an antichiral superfield, while \( \overline{\star} \)-product of chiral superfields is a chiral superfield. This means that we can write only antichiral superpotential of the form:

\[ \overline{W}_{(\star)} = \sum_n g_n \overline{\Phi}_n^\alpha = g_2 \overline{\Phi} \star \overline{\Phi} + g_3 \overline{\Phi} \star \overline{\Phi} \star \overline{\Phi} + \ldots \quad (19) \]
and its Hermitian conjugated chiral superpotential of the form:
\[ W(\bar{\Phi}) = \sum_n g_n \Phi^*_n = g_2 \Phi \Phi + g_3 \Phi \Phi \Phi + \ldots \] (20)

Moreover, ⋆-product (⋆̅-product) of antichiral (chiral) superfields is equivalent to the ordinary product, that is to say \( W(\Phi) W(\bar{\Phi}) = W(\Phi \bar{\Phi}) \), and thus the superpotential of the undeformed Wess-Zumino model is consistent with deformed supersymmetry under consideration. Also it is easy to see, the ⋆-product of chiral and antichiral superfields (the kinetic term of the Wess-Zumino model) is not deformed
\[ \Phi \Phi = \Phi \Phi, \] (21)
because of \( D_\alpha \Phi = 0 \). Thus we conclude that the ordinary undeformed Wess-Zumino model is fully compatible with deformed supersymmetry!\(^3\)

4 Vector superfields

Consider now SUSY gauge theory with arbitrary gauge group. The element of the SUSY gauge group is given by a chiral superfield \( \exp (i\Lambda) \), where \( \Lambda = \Lambda^a T^a \) are the chiral superfields \((\bar{D}_a \Lambda^a = 0)\) and \( T^a \) are the generators of the group. The antichiral element of the SUSY gauge group then is \( \exp (-i\bar{\Lambda}) \), \((D_a \bar{\Lambda}^a = 0)\). Once again, in order to preserve chirality of those superfields we should consider the following deformations: \( \exp \bar{\Lambda} \) and \( \exp (-i\bar{\Lambda}) \), which are equivalent to the corresponding undeformed superfields. This means that contrary to the case of ordinary noncommutative gauge theories the algebra of gauge group is not deformed and we have no restrictions on possible gauge groups and their representations. Actually, we will see momentarily that like the Wess-Zumino model discussed in the previous section, the ordinary undeformed gauge theories are also compatible with deformed SUSY.

The gauge fields are residing in a vector superfield \( V \), which is a Hermitian matrix, \( V^\dagger = V \). The gauge symmetry acts
\[ e^V_\dagger \rightarrow \exp \bar{\Phi} (i\Lambda) \exp \Phi \exp (i\Lambda) = \exp (-i\bar{\Lambda}) e^V \exp (i\Lambda), \] (22)
i.e. as in the case of undeformed SUSY. This means that we can use ordinary Wess-Zumino gauge, where
\[ V(x, \theta, \bar{\theta}) = -\sigma^\mu \bar{\theta} A_\mu (x) + i\theta \bar{\theta} \bar{\Phi} \lambda (x) - i\theta \bar{\theta} \bar{\Phi} \lambda (x) + \frac{1}{2} \theta \bar{\theta} \bar{\Phi} \lambda (x). \] (23)

\(^3\)In [18], it has been pointed out that the deformed N=1 supersymmetry does allow to write D-terms of the ⋆-product of chiral superfields, \( \bar{D}^2 \sum_n g_n \Phi^*_n \) (along with their Hermitian conjugate terms), which give a deformation of the Wess-Zumino model involving derivatives of the component fields. Obviously, such kind of terms in the classical (undeformed) limit, \( C \rightarrow 0 \), vanish.
Correspondingly we define chiral and antichiral field strength superfields:

\[ W^{(\overline{\mathcal{C}})}_\alpha = -\frac{1}{4} D^D e^V \overline{\mathcal{C}} D_\alpha e^V, \]
\[ W^{(*)}_\alpha = -\frac{1}{4} D^D e^V \star D_\alpha e^V. \]  \hspace{1cm} (24)

Using (22), one sees that they transform under the gauge transformations as in the usual case

\[ W^{(\overline{\mathcal{C}})}_\alpha \rightarrow \exp (-i\Lambda) W^{(\overline{\mathcal{C}})}_\alpha \exp (i\Lambda) \]
\[ W^{(*)}_\alpha \rightarrow \exp (-i\tilde{\Lambda}) W^{(*)}_\alpha \exp (i\tilde{\Lambda}). \]  \hspace{1cm} (25)

In fact, it is easy to see that field strength superfields (24) are not deformed. Indeed,

\[ W^{(\overline{\mathcal{C}})}_\alpha = -\frac{1}{4} D^D e^V \overline{\mathcal{C}} D_\alpha e^V = -\frac{1}{4} D^D \left( D_\alpha V - \frac{1}{2} [D_\alpha V, V] \right) = -\frac{1}{4} D^D \left( D_\alpha V - \frac{1}{2} [D_\alpha V, V] \right) - \frac{1}{8} D^D D^\beta D^\gamma D_\alpha e^V \overline{\mathcal{C}}^\beta \gamma \left[ D_\gamma V \overline{\mathcal{C}}^\beta \gamma D_\alpha V + V \overline{\mathcal{C}}^\beta \gamma D_\alpha V \right] + \frac{\det \overline{\mathcal{C}}}{16} \left[ \overline{\mathcal{C}}^\beta \gamma D_\alpha V \overline{\mathcal{C}}^\beta \gamma V - \overline{\mathcal{C}}^\beta \gamma V \overline{\mathcal{C}}^\beta \gamma D_\alpha V \right] \]
\[ = W_\alpha \]  \hspace{1cm} (26)

and similarly for \( \overline{W}^{(*)}_\alpha, \overline{W}^{(*)}_\alpha = \overline{W}_\alpha \). Then the invariant Lagrangian

\[ \mathcal{L}_{\star SYM} = \int d^2 \theta \frac{1}{4} W \overline{\mathcal{C}} W + \int d^2 \theta \frac{1}{4} \overline{W} \star \overline{W} \]  \hspace{1cm} (27)

is indeed nothing but the Lagrangian of ordinary super-Yang-Mills theory!

5 Conclusions

We have discussed the deformation of N=1 four dimensional Minkowski superspace assuming that the odd superspace coordinates \( \theta^a \) do not anticommute. As it has been recently shown in [14], in Euclidean space this deformation can be described by a star product which respects only \( N=\frac{1}{2} \) supersymmetry. In Minkowski space-time this deformation can be described by \( N=1 \) supersymmetric \( \star \)-product which is non-Hermitian. Hermiticity of Lagrangians ensures that this \( \star \)-deformation and its conjugated \( \overline{\star} \)-deformation are physically equivalent. The (anti)commutation relations among other superspace coordinates are also modified (see, (16, 17)). In particular, ordinary space-time coordinates do not commute as well.

We have defined chiral and vector superfields and discussed their interactions. Although in general one is able to write corrections to the classical (undeformed) \( N=1 \) Wess-Zumino model, but it is amusing that the ordinary Wess-Zumino model alone is fully compatible with deformed supersymmetry. We have demonstrated that the algebra of gauge groups are not deformed and hence the ordinary super-Yang-Mills theories are also consistent with the deformation of supersymmetry we have considered.
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