GOLDSTONE-TYPE SUPERFIELDS AND PARTIAL
SPONTANEOUS BREAKING OF D=3, N=2
SUPERSYMMETRY

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Abstract

We consider the modified superfield constraints with constant terms for the $D=3$, $N=2$ Goldstone-Maxwell gauge multiplet which contains Goldstone fermions, real scalar and vector fields. The partial spontaneous breaking $N=2 \to N=1$ is possible for the non-minimal self-interaction of this modified gauge superfield including the linear Fayet-Iliopoulos term. The dual description of the partial breaking in the model of the self-interacting Goldstone chiral superfield is also discussed.

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1 Goldstone-Maxwell superfield

Models with the partial spontaneous breaking of the global $D=3$, $N=2$ supersymmetry have been constructed using the topologically non-trivial classical solutions preserving the one half of supercharges [1] and also in the method of nonlinear realizations of supersymmetries using superfields of the unbroken $N=1$ supersymmetry [2]. These models describe interactions of the Goldstone fermions with the complex scalar field of the supermembrane or with the real scalar and vector fields of the $D2$-brane.

The standard linear supermultiplets (standard superfields) are not convenient for a description of the partial spontaneous breaking of the extended global supersymmetries ($PSBGS$) when the invariance with respect to the part of supercharges remains unbroken. The partial breaking of the $N=2$ supersymmetry means a degeneracy of the matrix of vacuum supersymmetric transformations of two real fermions in some supermultiplet, however, this is impossible for the standard structure of auxiliary components in the vector or chiral superfields. Nevertheless, the linear $N=2$ transformations with the partial breaking can be constructed in terms of the $N=1$ superfields [2]. We shall show that these representations with the linear Goldstone ($LG$) fermions correspond to the modified constraints for the Goldstone-type superfields in the $N=2$ superspace.

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The Goldstone-Maxwell chiral superfield $W$ in the $D=4, N=2$ superspace satisfies the modified superfield 2-nd order constraints \[3, 4\]. In comparison to the original constraints of the $N=2$ vector multiplet \[5\], the deformed constraints contain the \textit{constant} terms which guarantee the appearance of the unusual constant imaginary part of the isovector auxiliary component and the Goldstone fermion component in the gauge superfield. This abelian gauge model has been used to break spontaneously $D=4, N=2$ supersymmetry to its $N=1$ subgroup.

The more early example of the Goldstone-type constraint has been considered in the model with the partial breaking of the $D=1, N=4$ supersymmetry \[6\]. Thus, these constraints introduce a new type of the supersymmetry representations with the $LG$-fermions. In distinction with the Goldstone fermions of the nonlinear realizations which transforms linearly only in the unbroken supersymmetry, the $LG$-fermions have their partners in the supermultiplets of the whole supersymmetry. The nonlinear deformation of the standard constraints is also possible \[7\], however, we shall discuss only constant terms in the modified constraints which are connected with the spontaneous breaking of supersymmetries. It will be shown that the models with the $LG$ vector multiplet and the corresponding dual scalar multiplet solve the problem of the partial spontaneous breaking of the $D=3, N=2$ supersymmetry.

In this section we discuss the prepotential solution for the $LGM$ supermultiplet which contains additional terms manifestly depending on the spinor coordinates and some complex constants playing the role of moduli in the vacuum state of the theory together with the constant of the Fayet-Iliopoulos ($FI$) term. Using this representation in the non-minimal gauge action one can obtain the constant vacuum solutions with the partial spontaneous breaking of the $D=3, N=2$ supersymmetry. Note that the supersymmetry algebra is modified on the $LGM$ prepotential $V$ by analogy with the similar modified transformations of the 4D gauge fields or prepotentials in refs.\[3, 4\].

The sect.2 is devoted to the description of $PSBGS$ in the interaction of the $LG$-chiral superfield which is dual to the interaction of the $LGM$ superfield. This manifestly supersymmetric action depends on the sum of the chiral and antichiral superfields and some constant term bilinear in the spinor coordinates. The non-usual transformation of the basic $LG$-chiral superfield satisfies the supersymmetry algebra with the central-charge term.

The $N=1$ supermembrane and $D2$-brane actions \[2\] can be analysed in our approach using the decompositions of $N=2$ superfields in the 2-nd spinor coordinate $\theta^a_2$. We consider the $N=1$ components of the extended superfields and the covariant conditions which allow us to express the additional degrees of freedom in terms of the Goldstone superfields.

The coordinates of the full $D=3, N=2$ superspace are
\[
z = (x^{\alpha\beta}, \theta^a, \bar{\theta}^\alpha),
\]
where $\alpha, \beta$ are the spinor indices of the group $SL(2,R)$. The spinor representation of the coordinate is connected with the vector representation via the $3D\gamma$-matrices $x^{\alpha\beta}=(1/2)x^m(\gamma_m)^{\alpha\beta}$. The spinor derivatives in this superspace have the following form:
\[
D_\alpha = D_\alpha + \frac{i}{2}\bar{\theta}_a Z , \quad D_\alpha = \partial_\alpha + \frac{i}{2}\bar{\theta}^a \partial_\alpha \beta , \\
\bar{D}_\alpha = \bar{D}_\alpha - \frac{i}{2}\theta_\alpha Z , \quad \bar{D}_\alpha = \bar{\partial}_\alpha + \frac{i}{2}\theta^a \partial_\alpha \beta,
\]
where $Z$ is the real central charge, and $D_\alpha$ and $\bar{D}_\alpha$ are the spinor derivatives without the central charge.
The \( N=2 \) supersymmetry algebra is covariant with respect to the \( U_R(1) \) transformations of the spinor coordinates.

We shall consider the following notation for the bilinear combinations of spinor coordinates and differential operators:

\[
(\theta)^2 = \frac{1}{2} \theta^\alpha \bar{\theta}^\alpha , \quad (\bar{\theta})^2 = \frac{1}{2} \bar{\theta}^\alpha \bar{\theta}^\alpha ,
\]

\[
(\theta \bar{\theta}) = \frac{1}{2} \bar{\theta}^\alpha \theta^\alpha , \quad \Theta^{\alpha \beta} = \frac{1}{2} [\theta^\alpha \bar{\theta}^\beta + \alpha \leftrightarrow \beta] ,
\]

\[
(D)^2 = \frac{1}{2} D^\alpha D^\alpha , \quad (\bar{D})^2 = \frac{1}{2} \bar{D}_\alpha \bar{D}^\alpha ,
\]

\[
(D \bar{D}) = \frac{1}{2} D^\alpha \bar{D}_\alpha , \quad D_{\alpha \beta} = \frac{1}{2} ([D^\alpha , \bar{D}_\beta] + \alpha \leftrightarrow \beta) .
\]

The complex chiral coordinates can be constructed by the analogy with \( D=4 \)

\[
\zeta = (x^{\alpha \beta}, \theta^\alpha) , \quad x^{\alpha \beta} = x^{\alpha \beta} + i \Theta^{\alpha \beta} .
\]

It is convenient to use the following rules of conjugation for any operators [7]:

\[
(XY)^\dagger = Y^\dagger X^\dagger , \quad [X,Y]^\dagger = -(-1)^{p(X)p(Y)} [X^\dagger , Y^\dagger] ,
\]

where \([X,Y]\) is the graded commutator and \(p(X) = \pm 1\) is the \( Z_2 \)-parity.

It is possible to introduce the real \( N=2 \) spinor coordinates \( \theta^\alpha_k = (\theta^\alpha_k)^\dagger \)

\[
\theta^\alpha_1 = \frac{1}{\sqrt{2}} (\theta^\alpha + \bar{\theta}^\alpha) , \quad \theta^\alpha_2 = \frac{i}{\sqrt{2}} (\bar{\theta}^\alpha - \theta^\alpha) ,
\]

\[
(\theta \bar{\theta}) = \frac{1}{2} [((\theta_1 \theta_1) + (\theta_2 \theta_2)] , \quad (\theta_i \theta_k)^\dagger = -((\theta_i \theta_k) ,
\]

and the corresponding real spinor derivatives

\[
D^1_\alpha = D^1_\alpha + \frac{1}{2} \theta_2 \alpha Z , \quad D^2_\alpha = D^2_\alpha - \frac{1}{2} \theta_1 \alpha Z ,
\]

\[
D^1_\alpha = \frac{1}{\sqrt{2}} (D^\alpha + \bar{D}_\alpha) , \quad D^2_\alpha = \frac{i}{\sqrt{2}} (D^\alpha - \bar{D}_\alpha) .
\]

The \( D=3, N=2 \) gauge theory [8, 9] is analogous to the well-known \( D=4, N=1 \) gauge theory, although the three-dimensional case has some interesting peculiarities which are connected with the existence of the topological mass term and duality between the 3D-vector and chiral multiplets. We shall consider the basic superspace with \( Z=0 \).

The abelian \( U(1) \)-gauge prepotential \( V(z) \) possesses the gauge transformation \( \delta V = \Lambda + \bar{\Lambda} \) where \( \Lambda \) is the chiral parameter.

The \( D=3, N=2 \) vector multiplet is described by the real linear superfield

\[
W(V) = i (D \bar{D}) V
\]

satisfying the basic constraints \((D)^2 W = (\bar{D})^2 W = 0\).

The components of \( W(V) \) are the real scalar \( \varphi \), the field-strength of the gauge field \( F_{\alpha \beta}(A) \), the real auxiliary component \( G \) and the spinor fields \( \lambda^\alpha \) and \( \bar{\lambda}^\alpha \).
The low-energy effective action of the 3D vector multiplet describes a non-minimal interaction of the real scalar field with the fermion and gauge fields. For the \( U(1) \) gauge superfield \( V \) this action has the following general form:

\[
S(W) = -\frac{1}{2} \int d^7 z H(W) , \quad \tau(W) = H''(W) > 0 , \quad (1.14)
\]

where \( H(W) \) is the real convex function of \( W \).

Let us consider the spontaneous breaking of supersymmetry in the non-minimal gauge model (1.14) with the additional linear FI-term

\[
S_{FI} = \frac{1}{2} \xi \int d^7 z V , \quad (1.15)
\]

where \( \xi \) is a constant of the dimension \(-1\). Varying the superfield \( V \) one can derive the corresponding superfield equation of motion. We shall study the constant solutions of this equation using the following vacuum ansatz:

\[
V_0 = 2i(\bar{\theta}\theta)a - 2(\theta)^2(\bar{\theta})^2G , \quad W_0 = a + 2i(\theta\bar{\theta})G , \quad (1.16)
\]

where \( a \) and \( G \) are constants. The non-trivial solution \( G_0 \neq 0 \) is possible for the quadratic function \( H \) only.

It is useful to consider the real spinors \( \lambda_k^0 \) and the real spinor parameters of the \( N=2 \) supersymmetry by analogy with (1.9). It is clear that the constant solution \( G_0 \neq 0 \) can only break spontaneously both supersymmetries.

Consider the following deformation of the linearity constraints which defines the \( LGM \) superfield:

\[
(D)^2 \hat{W} = C , \quad (\bar{D})^2 \hat{\bar{W}} = \bar{C} , \quad (1.17)
\]

where \( C \) and \( \bar{C} \) are some constants. These relations are manifestly supersymmetric, however, they break the \( U_R(1) \) invariance.

The gauge-prepotential solution of these constraints can be constructed by analogy with Eq.(1.13)

\[
\hat{W} = i(DD)V + (\theta)^2C + (\bar{\theta})^2\bar{C} . \quad (1.18)
\]

This superfield contains new constant auxiliary structures which change radically the matrix of the vacuum fermion transformations \( \delta k \lambda_k^0 \). It is evident that the \( PSBGS \)-condition corresponds to the degeneracy of these transformations

\[
CC - G_0^2 = 0 . \quad (1.19)
\]

In this case one can choose the single real Goldstone spinor field as some linear combination of \( \lambda_k^0 \).

The action of the \( LGM \)-superfield (1.18) has the following form:

\[
\hat{S}(V) = -\frac{1}{2} \int d^7 z [H(\hat{W}) - \xi V] \quad (1.20)
\]

and depends on three constants \( \xi, C \) and \( \bar{C} \).

The non-derivative terms in the component Lagrangian produce the following scalar potential

\[
\mathcal{V}(\varphi) = \frac{1}{2} [ |C|^2 \tau(\varphi) + \xi^2 \tau^{-1}(\varphi) ] . \quad (1.21)
\]
The PSBGS solution (1.13) arises for the non-trivial interaction $\tau'(a) \neq 0$. This solution determines the minimum point $a_0$ of this model

$$\tau(a_0) = \frac{|\xi|}{|C|}.$$  

(1.22)

The vacuum auxiliary field can be calculated in the point $a_0$

$$G_0 = \frac{\xi}{\tau(a_0)} = \pm |C|.$$  

(1.23)

Using the $U_R(1)$ transformation one can choose the pure imaginary constant $C \rightarrow c = i|c|$ (without the loss of generality) then $G_0 = -ic = |c|$.

This choice corresponds to the following decomposition of the LGM-superfield (1.18)

$$\hat{W} = W_s(V_s) + 2ic|\theta_2\theta_2|,$$

$$W_s(V_s) = i\alpha(\alpha^2D^2_1 + D^2_2) V_s.$$  

(1.24)

(1.25)

where $V_s$ is the shifted LGM-prepotential which has the vanishing vacuum solution for the auxiliary component. It is evident that this representation breaks spontaneously the 2-nd supersymmetry only.

It should be stressed that the shifted quantities $W_s$ and $V_s$ are not standard superfields

$$\delta_\epsilon W_s = i\epsilon_k^kQ^k_\alpha \hat{W} = -2ic|\epsilon^2_2\theta_2 + i\epsilon_k^kQ^k_\alpha W_s,$$

$$\delta_\epsilon V_s = 2c|\epsilon^2_2\theta_2(\theta_1^2\theta_1) + i\epsilon_k^kQ^k_\alpha V_s.$$  

(1.26)

(1.27)

The supersymmetry algebra of the $V_s$-transformations is essentially modified by the analogy with the transformations of the prepotentials in refs. [4, 11].

It should be remarked that the minimal interaction of the charged chiral superfields with the LGM-prepotential $V_s$ breaks the supersymmetry. The analogous problem of the LGM interaction with the charged matter appears also in the PSBGS model with $D=4$, $N=2$ supersymmetry [4].

2 Goldstone chiral superfield

The 3D linear multiplet is dual to the chiral multiplet $\phi$. The Legendre transform describing this duality is

$$S[B, \Phi] = -\frac{1}{2} \int d^7z [H(B) - \Phi B],$$  

(2.1)

where $B$ is the real unconstrained superfield and $\Phi = \phi + \bar{\phi}$. Varying the Lagrange multipliers $\phi$ and $\bar{\phi}$ one can obtain the linearity constraints for $B$.

We shall show that the spontaneous breaking of supersymmetry is possible for the non-trivial interaction of the LG chiral superfield which possesses the inhomogeneous supersymmetry transformation. Let us consider the dual picture for the PSBGS gauge model with the FI-term (1.20)

$$\hat{S}(B, \phi, \bar{\phi}) = -\frac{1}{2} \int d^7z [H(B) - B\Phi] - \frac{1}{2} [\bar{C} \int d^3x(D)^2\phi + c.c.],$$  

(2.2)
where the modified constrained $LG$ superfield is introduced
\[
\hat{\Phi} \equiv \phi + \bar{\phi} + 2i\xi(\theta\bar{\theta}) , \quad (D\bar{D})\hat{\Phi} = -i\xi .
\] (2.3)

Varying the chiral and antichiral Lagrange multipliers $\phi$ and $\bar{\phi}$ one can obtain the $LGM$-constraints (1.17) on the superfield $B$ and then pass to the gauge phase $B \rightarrow \hat{W}(V)$ where the $(\theta\bar{\theta})$-term in $\hat{\Phi}$ transforms to the $FI$-term.

The algebraic $B$-equation
\[
H'(B) \equiv f(B) = \hat{\Phi} ,
\] (2.4)
provides the transform to the ‘chiral’ phase
\[
B \rightarrow f^{-1}(\hat{\Phi}) \equiv \hat{B}(\hat{\Phi}) .
\] (2.5)

The transformed chiral action is
\[
\hat{S}(\hat{\Phi}) = -\frac{1}{2} \int d^7z \{ \hat{H}(\hat{\Phi}) + [C(\theta)^2 + c.c.]\hat{\Phi} \} ,
\] (2.6)
\[
\hat{H}(\hat{\Phi}) = H[\hat{B}(\hat{\Phi})] - \hat{\Phi}\hat{B}(\hat{\Phi})
\] (2.7)
The linear terms with $C$ and $\bar{C}$ break the $U_R(1)$-symmetry, however, this action is invariant with respect to the isometry transformation.

It should be underlined that the $LG$-superfield $\hat{\Phi}$ transforms homogeneously, while the supersymmetry transformation of the $LG$-chiral Lagrange multiplier $\phi$ contains the inhomogeneous term
\[
\delta_\epsilon \phi = -i\xi (\theta^a\bar{\epsilon}_a) + i\epsilon^a_k Q^k \phi .
\] (2.8)

Consider the $\theta$-decomposition of the $LG$-chiral superfield
\[
\phi = A(x_L) + \theta^a \psi_a(x_L) + (\theta)^2 F(x_L) ,
\] (2.9)
where $x_L$ is the coordinate of the chiral basis.

The Lie bracket of the modified supersymmetry transformation (2.8) contains the composite central charge parameter corresponding to the following action of the generator $Z$ on the chiral superfield:
\[
Z\phi = \xi , \quad (Z\bar{\phi} = -\xi) .
\] (2.10)
Thus, the Goldstone boson field $\text{Im} A(x)$ for the central-charge transformation appears in this model. It should be remarked that the isometry transformation in the chiral model without $PSBGS$ cannot be identified with the central charge.

The vacuum equations of motion for this model have the following form:
\[
\bar{F}\hat{\tau}(b) + \bar{C} = 0 , \quad b = A + \bar{A}
\] (2.11)
\[
(|F|^2 - \xi^2)\hat{\tau}'(b) = 0 , \quad \hat{\tau} = \hat{H}'' = -\tau^{-1} .
\] (2.12)

The scalar potential of this model depends on the one real scalar component only
\[
\mathcal{V}(b) = \frac{1}{2}[\xi^2\hat{\tau}(b) + |C|^2\hat{\tau}^{-1}(b)] .
\] (2.13)

The vacuum solution $|F_0|^2 = \xi^2$ corresponds to the degeneracy condition for the matrix of vacuum supersymmetry transformations. The choice $F_0 = i\xi$ breaks the 2-nd supersymmetry.
Thus, the non-trivial interaction of the $LG$-chiral superfield $\phi$ provides the partial spontaneous breaking of the $D=3$, $N=2$ supersymmetry. This phenomenon has been analysed also in the formalism of the $D=3$, $N=1$ Goldstone-type superfields [2].

Let us assume that the spinor coordinates $\theta_1^\alpha$ parameterize $N=1$ superspace, and the generators $Q_1^a$ form the corresponding subalgebra of the $N=2$ supersymmetry.

The chirality condition in the real basis

$$(D_1^a + iD_2^a)\phi = 0$$  \hspace{1cm} (2.14)

can be solved via the complex unrestricted $N=1$ superfield

$$\phi = \chi(x, \theta_1) + i\theta_2^a D_1^a \chi(x, \theta_1) + (\theta_2 \theta_2)(D^1 D^1) \chi(x, \theta_1) .$$  \hspace{1cm} (2.15)

The transformation (2.8) generates the corresponding transformation of the complex $N=1$ superfield:

$$\delta \chi = - \frac{1}{2} [\xi \theta_1^\alpha (\epsilon_{2\alpha} + i\epsilon_{1\alpha}) - i\epsilon_2^a D_1^a \chi + i\epsilon_1^a Q_1^a \chi] .$$  \hspace{1cm} (2.16)

Consider the $\theta_2$-decomposition of the basic superfield (2.3) of the chiral $PSBGS$ model

$$\hat{\Phi} = \chi + \bar{\chi} + i\theta_2^a D_1^a (\chi - \bar{\chi}) + (\theta_2 \theta_2)(D^1 D^1)(\chi + \bar{\chi}) + i[\epsilon(\theta_1 \theta_1) + (\theta_2 \theta_2)]$$

$$= \Sigma + \theta_2^a D_1^a \rho + (\theta_2 \theta_2)[(D^1 D^1) \Sigma + 2i\xi] ,$$  \hspace{1cm} (2.17)

where $\Sigma$ is the massive real $N=1$ superfield and $\rho$ is the real Goldstone superfield for the 2-nd supersymmetry.

The $N=1$ superfields with the analogous $N=2$ transformations have been proposed in ref.[2]. The authors of this work have shown that the additional superfield can be constructed in terms of the spinor derivative of the Goldstone superfield $\rho$ in order to built the supermembrane action. In our approach, the massive degrees of freedom can be removed using the covariant condition

$$\hat{\Phi}^2 = 0 ,$$  \hspace{1cm} (2.18)

which allows us to construct $\Sigma$ via $D_1^a \rho$ by analogy with the similar construction in the $D=4$, $N=2$ theory [12].

Let us analyse the $N=1$ decomposition of the gauge prepotential

$$V_s(x, \theta_1, \theta_2) = \kappa(x, \theta_1) + i\theta_2^a V_\alpha(x, \theta_1) + i(\theta_2 \theta_2) M(x, \theta_1)$$  \hspace{1cm} (2.19)

and the chiral gauge parameter $\Lambda = \exp(i\theta_2^a D_1^a) \lambda(x, \theta_1)$.

The gauge transformations of the $N=1$ components are

$$\delta_\lambda \kappa = \lambda + \bar{\lambda} , \hspace{1cm} \delta_\lambda V_\alpha = D_1^a (\lambda - \bar{\lambda}) .$$  \hspace{1cm} (2.20)

Thus, $\kappa$ is a pure gauge degree of freedom, $V_\alpha$ is the $N=1$ gauge superfield, and $M$ is the scalar $N=1$ component of the $N=2$ supermultiplet.

Consider the $N=1$ decomposition of the linear superfield (1.25)

$$W_s(V_s) = w + i\theta_2^a F_\alpha(V) - (\theta_2 \theta_2)(D^1 D^1) w ,$$  \hspace{1cm} (2.21)

where the gauge-invariant scalar and spinor superfields are defined

$$w = \frac{1}{2} [M + i(D^1 D^1) \kappa] ,$$  \hspace{1cm} (2.22)

$$F_\alpha(V) = i \frac{1}{2} (D^1 D^1) V_\alpha + \frac{1}{4} \theta_\alpha \beta V_\beta , \hspace{1cm} D^\alpha F_\alpha = 0 .$$  \hspace{1cm} (2.23)
The Goldstone transformation of $W_s$ produces the $\epsilon_2$-transformations of the $N=1$ superfields. The spinor superfield strength $F_\alpha$ is analogous to the Goldstone spinor superfield of ref.\cite{2}. It describes the Goldstone degree of freedom of the D2-brane, and the superfield $w$ corresponds to the massive degrees of freedom. Our construction introduces the $N=1$ gauge superfield $V_\alpha$ as the basic object of this model and allows us to study the modification of the supersymmetry algebra on the gauge fields of the D2-brane.

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