SUPERSYMMETRIC 3-BRANE FROM THE EXTENDED GOLDSTONE-MAXWELL SUPERMULTIPLET

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A new gauge supermultiplet of the N=2, D=4 supersymmetry with the complex central charge is proposed. Besides the global N=2, D=4 supersymmetry it possesses the symmetry under the central charge transformations and the special N=2 gauge invariance. It is shown that when one of the supersymmetries is spontaneously broken and the gauge is fixed this multiplet describes the standard action of the supersymmetric 3-brane.

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1. INTRODUCTION

The supersymmetric 3-brane is the most striking pattern of theory where the partial breaking of global supersymmetry (PBGS) was firstly revealed [1]. It was realized that in this system we deal with the Nambu-Goldstone mechanism which is responsible for the breaking of one half of the supersymmetries and where the Goldstone fermion belongs to a multiplets of the residual unbroken supersymmetry. As we know from [2] the most direct approach to such systems is the nonlinear realization method which gives the action of the super-3-brane in terms of the unconstrained N=1 chiral worldvolume Goldstone superfield φ. However, this method becomes complicated in the process of deriving of the action due to the absence of the direct procedure of constructing of the relevant linear representation of N=2, D=4 supersymmetry in terms of constrained Goldstone superfields. In [3, 5] there was proposed another approach in which the required Goldstone superfield is embedded into the linear N=2 tensor representation constrained by the suitable nilpotence relations. However, in order to get the supersymmetric 3-brane action in terms of the scalar Goldstone φ in this approach the additional procedure of the duality transformation is required.

In this contribution we would like to propose a new approach to this problem which is based on unknown up to now linear representation of N=2, D=4 supersymmetry. Besides the N=2, D=4 Goldstone-Maxwell superfield the latter involves the new gauge prepotential composed out of the two N=2, D=4 superfields with the different central charges and the N=2 gauge transformations. In contrast to the aforementioned supermultiplet the new one describes the linear representation of N=2, D=4 supersymmetry which the Goldstone superfield φ and the corresponding Lagrangian density L of the super-3-brane belong to. We will show that this approach gives the straightforward way of construction of the corresponding action provided the reasonable nonlinear realization of the supersymmetry considered is taken into account.

2. N=2 PREPOTENTIAL WITH CENTRAL CHARGE

2.1. EXTENDED GOLDSTONE-MAXWELL SUPERMULTIPLET

Let us consider the following central extension of N=2, D=4 Poincare superalgebra

$$\{Q_\alpha, \bar{Q}_{\dot\alpha}\} = -2i\partial_{\alpha\dot\alpha}, \{S_\alpha, \bar{S}_{\dot\alpha}\} = -2i\partial_{\alpha\dot\alpha}, \quad (1)$$
$$\{Q_\alpha, S_\beta\} = 2\varepsilon_{\alpha\beta} Z, \{\bar{Q}_\alpha, \bar{S}_\beta\} = -2\varepsilon_{\dot\alpha\dot\beta}\bar{Z}$$

with the other (anti)commutators vanishing. It is well known that this supersymmetry can be realized on the N=2 vector Goldstone supermultiplet

$$\delta\varphi = f + \eta^\alpha W_\alpha, \quad f = c - 2i\eta^\alpha\theta_\alpha, \quad (2)$$
$$\delta W_\alpha = -\frac{1}{2}\bar{D}^2\bar{\varphi}\eta_\alpha - 2i\partial_{\alpha\dot\alpha}\varphi\bar{\eta}^{\dot\alpha},$$

where φ and W_α are chosen to be chiral

$$\bar{D}_{\dot\alpha}\varphi = 0, \quad \bar{D}_{\dot\alpha} W_\alpha = 0,$$

and constrained by the reality condition

$$D^\alpha W_\alpha + \bar{D}_{\dot\alpha}\bar{W}^{\dot\alpha} = 0.$$

Note that we use the following conjugation rules for the vector and spinor derivatives [6] $(\partial_\alpha)^+ = -\bar{\partial}_{\dot\alpha}$, $(\partial_{\alpha\dot\alpha})^+ = \partial_{\alpha\dot\alpha}$, $(D_\alpha)^+ = \bar{D}_{\dot\alpha}$, $(D_\alpha\bar{D}_{\dot\alpha})^+ = -\bar{D}_{\dot\alpha}D_\alpha$, $(D^\alpha D_\alpha)^+ = \bar{D}_{\dot\alpha}\bar{D}^{\dot\alpha}$.

Recall that the latter can be solved in terms of the corresponding gauge prepotential

$$W_\alpha = -i/4\,\bar{D}^2 D_\alpha L, \quad \bar{W}_{\dot\alpha} = (W_\alpha)^+ = i/4\, D^2\bar{D}_{\dot\alpha} L.$$

The superfields φ, W_α can be embedded into the N=2 D=4 superfield

$$W = \varphi - \omega^\alpha W_\alpha - 1/4\,\omega^2\bar{D}^2\bar{\varphi} - i\omega^\alpha\bar{\omega}^{\dot\alpha}\partial_{\alpha\dot\alpha}\varphi +$$
$$+ i/2\,\omega^2\bar{\omega}_{\dot\alpha}\partial^{\alpha\dot\alpha}W_\alpha - 1/8\,\omega^2\bar{\omega}^2\partial^{\alpha\dot\alpha}\partial_{\alpha\dot\alpha}\varphi,$$

restricted by the chirality conditions

$$\bar{D}^i_{\dot\alpha} W = D^i_\alpha \bar{W} = 0, \quad (3)$$
$$D^i_\alpha = \partial/\partial\theta^\alpha_i - i\bar{\theta}^{\dot\alpha i}\partial_{\alpha\dot\alpha},$$

and the Bianchi identity

$$D^{ik} W = \bar{D}^{ik}\bar{W}, \quad (4)$$
$$D^{ik} \equiv D^{i\alpha} D^k_\alpha, \quad \bar{D}^{ik} \equiv \bar{D}^i_{\dot\alpha}\bar{D}^{k\dot\alpha}.$$

It will be useful for us to keep the following notations for the Grassmann coordinates and spinor covariant derivatives

$$\theta_1^\alpha = \theta^\alpha, \theta_2^\alpha = \omega^\alpha, \; D_\alpha^1 = D_\alpha, \; D_\alpha^2 = D_\alpha^{(\omega)}.$$

In these notations the constraints (3), (4) can be rewritten as

$$\bar{D}_{\dot\alpha} W = \bar{D}_{\dot\alpha}^{(\omega)} W = 0,$$
$$(D^{(\omega)})^2 W = \bar{D}^2 \bar{W}, \qquad (5)$$
$$D^{(\omega)\alpha} D_\alpha W = -\bar{D}_{\dot\alpha}^{(\omega)} \bar{D}^{\dot\alpha} \bar{W}.$$

In accordance to (2) the superfield W transforms as

$$\delta W = f + (\eta^\alpha S_\alpha + \text{h.c.})W, \qquad (6)$$

where S_α and their hermitian conjugate are the generators of 2-d supersymmetry similar to those in (1) but without the central charge contributions. From the Eq. (6) it follows that the S-supersymmetry is spontaneously broken and the N=1 superfield $D_\alpha\varphi$ can be treated as the corresponding Goldstone fermion. We remind, that, for example, in the case of the "space filling" D3-brane the analogous consideration, but without the central charge, amounts to the absolutely different result. It was found there that the related Goldstone fermion coincides with the component W_α [4,10]. Note, by the way, that in this case there exists the transparent explanation of the origin of the Nambu-Goldstone mechanism. It follows from the showing up the superfields that develop the nonzero vacuum expectation value among the components of W [7,8,10]. There is nothing like that in our case. So, first of all, we try to make a relevant extension of our supermultiplet to provide the conditions for the partial breaking of the N=2, D=4 supersymmetry of the pattern required in the case of theory with central charge. To this end let us add the set of two new N=2 superfields to the Goldstone-Maxwell superfield W and postulate that these superfields are connected with each other as follows

$$\bar{D}_{\dot\alpha} \Phi = \bar{D}_{\dot\alpha}^{(\omega)} \Phi = 0,$$
$$(D^{(\omega)})^2 \Phi = -\bar{D}^2 \Lambda, \qquad (7)$$
$$D_\alpha W = i/4\, (\bar{D}^{(\omega)})^2 D_\alpha^{(\omega)} \Lambda + 2i\beta\,\omega_\alpha.$$

From the dimensionality reasons we see that the new superfields have the equal dimensions $[\Phi]=[\Lambda]=cm^2$. The dimensionless constant β in the last constraint in (7) indicates that the v.e.v. of the superfield Λ is nonzero

$$\Lambda_{vac} = \beta\,\omega^2 \bar\omega^2.$$

It yields the PBGS of the pattern required. To understand that the extension at hand really gives the result desired it is very important to realize that the constraints (5) and (7) are covariant under the following S-supersymmetry transformations

$$\delta\Phi = fW + (\eta^\alpha S_\alpha + \text{h.c})\Phi,$$
$$\delta\Lambda = f\bar{W} + \bar{f}W + (\eta^\alpha S_\alpha + \text{h.c})\Lambda. \qquad (8)$$

Together with the Eq. (6) they form the full set of the transformations which display the nontrivial dependence of the supermultiplet W, Φ and Λ on the central charge. Thanks to this fact we are ready to go on to the central part of our investigation. Namely, we are going to show that the constraints (5) and (7) give an exhaustive explanation of PBGS stipulated by the Nambu-Goldstone mechanism. To get it we substitute the general Grassmann decompositions of the superfields Φ and Λ on their N=1 components into the constraints (5) and (7). The general solution of these constraints compatible with the transformations (6), (8) reads as follows

$$\Phi = F + \omega^\alpha \Psi_\alpha + 1/4\,\omega^2 \bar{D}^2 L - i\omega^\alpha \bar\omega^{\dot\alpha} \partial_{\alpha\dot\alpha} F -$$
$$- i/2\,\omega^2 \bar\omega_{\dot\alpha} \partial^{\dot\alpha\alpha} \Psi_\alpha - 1/8\,\omega^2 \bar\omega^2 \partial^2 F,$$
$$\Lambda = L + \omega^\alpha Z_\alpha + \bar\omega_{\dot\alpha} \bar{Z}^{\dot\alpha} -$$
$$-1/4\,\omega^2 \bar{D}^2 \bar\varphi - 1/4\,\bar\omega^2 D^2 \varphi -$$
$$- 1/2\,\omega^\alpha \bar\omega^{\dot\alpha} [D_\alpha, \bar{D}_{\dot\alpha}] L +$$
$$+ 1/2\,\omega^2 \bar\omega_\alpha (i\partial_\alpha{}^{\dot\alpha} Z^\alpha - 2i\bar{D}^{\dot\alpha} \bar\varphi) +$$
$$+ 1/4\,\omega^2 \bar\omega^2 (4\beta - 1/2\,\partial^2 L + iD^\alpha W_\alpha), \qquad (9)$$
$$\partial^2 \equiv \partial^{\alpha\dot\alpha} \partial_{\alpha\dot\alpha}.$$

In these decompositions the N=1 superfields F, Ψ_α, Z_α are chiral and satisfy the "duality" constraints

$$D^2 \Psi_\alpha = 4i\,\partial_{\alpha\dot\alpha} \bar{Z}^{\dot\alpha}, \; \bar{D}^2 \bar{Z}^{\dot\alpha} = 4i\,\partial^{\alpha\dot\alpha} \Psi_\alpha, \qquad (10)$$

while the real component L remains unconstrained. Now it is very instructive to enumerate the main results extracted from the solutions (9), (10).

1. Among the spinor components of the superfield Λ it appears one $2iD_\alpha\varphi - i\partial_{\alpha\dot\alpha} Z^{\dot\alpha}$ which acquires an inhomogeneous shift $\sim\beta$ in its transformation law proportional to the nonzero v.e.v. of the oldest component $N_{vac} = \beta$. Therefore, this component can be considered as the Goldstone fermion related with spontaneous breakdown of the S-supersymmetry while the Q-supersymmetry still unbroken.
2. Opposed to the vector multiplet the new one includes the prepotential L among its components.
3. At last there exists one remarkable feature of this extension procedure. It turns out that the auxiliary components of the new multiplet F and Ψ_α can be transformed away owing to the gauge invariance of the constraints (7).

2.2. GAUGE SYMMETRY AND THE ACTION

It is not hard to check that the constraints (7) are not changed under the action of following gauge transformations

$$\delta_g W = 0, \; \delta_g \Phi = X, \; \delta_g \Lambda = X + \bar{X}, \qquad (11)$$

where X is the chiral N=2 superfield constrained by the conditions

$$\bar{D}_{\dot\alpha} X = \bar{D}_{\dot\alpha}^{(\omega)} X = 0, \; (D^{(\omega)})^2 X = -\bar{D}^2 \bar{X}. \qquad (12)$$

It is obvious that these transformations are nothing but the N=2 generalization of the N=1 gauge freedom of the prepotential L. It gives the possibility of elimination of the two lower components of the superfield Φ. In this gauge one gets

$$\Phi = 1/4\,\omega^2 \bar{D}^2 L,$$

$$\Lambda_g = L - 1/2\,\omega^{\alpha}\bar{\omega}^{\dot\alpha}[D_\alpha, \bar{D}_{\dot\alpha}]L + \qquad (13)$$
$$+ i\bar{\omega}^2\omega^\alpha D_\alpha\varphi - i\omega^2\bar{\omega}_{\dot\alpha}\bar{D}^{\dot\alpha}\bar{\varphi} +$$
$$+ 1/4\,\omega^2\bar{\omega}^2(4\beta - 1/2\,\partial^2 L + iD^\alpha W_\alpha)\,.$$

Thus, we see that superfield Λ actually describes the Goldstone degrees of freedom of the spontaneously broken hidden gauge symmetry defined by the transformations (11). The gauge (13) is easily recognized as the *unitary* one. Note, that this gauge is very useful for the understanding of the dynamical meaning of the superfield Λ. It gives the simple form of the action of the super-3-brane

$$S = -1/2\int d^4x\, d^2\theta\, d^2\omega\, \Phi + \text{h.c.} = \int d^4x\, d^4\theta\, \Lambda\,, \qquad (14)$$

compatible with the all aforementioned symmetries. To be more precise, let's return to the transformations (11). One can check that the action (14) is invariant under the transformations (11) up to the surface terms, because, e.g. having been replaced by X the integrand with Λ becomes the x-derivative owing to the constraint (12). On the other hand, passing to the gauge (13) and performing the Grassmann integration via the variables ω we get the form

$$S = \int d^4x\, d^4\theta\, L\,. \qquad (15)$$

This action coincides explicitly with that one given in [3]. Thus, we can reduce our problem to the question of deriving of the nonlinear realization in which the Lagrangian density L could be expressed completely through the scalar Goldstone superfield φ and its covariant derivatives only.

3. NONLINEAR REALIZATION AND FURTHER CONSTRAINTS

3.1. EQUIVALENCE CONDITIONS

To deduce this expression let us shortly remind the algorithmic procedure of constructing the relevant nonlinear realization. Starting from the coset representation of the corresponding group element **g** one can define the transformation properties of the N=1, D=4 Goldstone superfields ψ_α and $\tilde\varphi$ [2]

$$\tilde{x}^{\alpha\dot\alpha\,\prime} = \tilde{x}^{\alpha\dot\alpha} + 2i\left(\eta^\alpha\bar\psi^{\dot\alpha}(\tilde{x},\theta,\bar\theta) - \psi^\alpha\bar\eta^{\dot\alpha}(\tilde{x},\theta,\bar\theta)\right),$$
$$\psi^{\alpha\,\prime}(\tilde{x}',\theta,\bar\theta) = \eta^\alpha + \psi^\alpha(\tilde{x},\theta,\bar\theta)\,,$$
$$\varphi'(\tilde{x}',\theta,\bar\theta) = f + \varphi(\tilde{x},\theta,\bar\theta)\,,$$

related to the generators of spontaneously broken symmetries S_α and Z[1]. These superfields can be covariantly constrained [2]

$$\psi_\alpha = i/2\,\nabla_\alpha\tilde\varphi\,,\ \bar\nabla_{\dot\alpha}\tilde\varphi = 0\,, \qquad (16)$$

where the appropriated covariant derivatives are introduced [11]

$$\nabla_{\alpha\dot\alpha} = T^{-1}{}^{\beta\dot\beta}_{\alpha\dot\alpha}\partial_{\beta\dot\beta}\,,$$
$$\nabla_\alpha = D_\alpha - i\left(\bar\psi^{\dot\beta}D_\alpha\psi^\beta + \psi^\beta D_\alpha\bar\psi^{\dot\beta}\right)\nabla_{\beta\dot\beta}\,,$$
$$\bar\nabla_{\dot\alpha} = \bar D_{\dot\alpha} - i\left(\bar\psi^{\dot\beta}\bar D_{\dot\alpha}\psi^\beta + \psi^\beta\bar D_{\dot\alpha}\bar\psi^{\dot\beta}\right)\nabla_{\beta\dot\beta}\,,$$

$$T^{\beta\dot\beta}_{\alpha\dot\alpha} = \delta^\beta_\alpha\delta^{\dot\beta}_{\dot\alpha} - i\psi^\beta\partial_{\alpha\dot\alpha}\bar\psi^{\dot\beta} - i\bar\psi^{\dot\beta}\partial_{\alpha\dot\alpha}\psi^\beta\,.$$

To be able to use this approach in the present framework of linear realization one may take advantage of the general procedure developed in [10, 12, 14]. Let us perform the restriction of the third constraint in (7) onto the hypersurface

$$x^{\alpha\dot\alpha} = \tilde{x}^{\alpha\dot\alpha}\,,\ \omega^\alpha = \psi^\alpha(\tilde{x},\theta,\bar\theta)\,. \qquad (17)$$

This is the key point in revealing the straightforward relationship between two kinds of realizations of supersymmetry – linear and nonlinear, because, in accordance with the general features of the underlying transformation laws, both of the superfields involving into the aforementioned constraint $D_\alpha W - 2i\beta\omega_\alpha$ and $(\bar D^{(\omega)})^2 D^{(\omega)}_\alpha\Lambda$ transform *independently* when restricted onto the hypersurface (17), for example

$$\delta\tilde\Sigma_\alpha = -i(\eta^\beta\bar\psi^{\dot\beta} - \bar\eta^{\dot\beta}\psi^\beta)\partial_{\beta\dot\beta}\tilde\Sigma_\alpha\,,$$
$$\tilde\Sigma_\alpha = (\bar D^{(\omega)})^2 D^{(\omega)}_\alpha\Lambda\Big|_{x=\tilde{x},\,\omega=\psi}\,.$$

Hence, one can impose the following covariant constraint

$$2i\beta\omega_\alpha = D_\alpha W\big|_{x=\tilde{x},\,\omega=\psi}\,. \qquad (18)$$

Eq. (18) establishes the sought for equivalent relationship between the Goldstone fermion of the nonlinear realization ψ_α we started with and its linear counterpart $D_\alpha\varphi$. Now, it is not hard to verify that the constraint (18) recovers the IHE-solution represented in (16) when the new nonlinear realization constraints are imposed

$$\tilde\varphi = W\big|_{x=\tilde{x},\,\omega=\psi}\,,\ D^{(\omega)}_\alpha W\big|_{x=\tilde{x},\,\omega=\psi} = 0\,. \qquad (19)$$

Resolving the Eq. (18) (with β=1) with respect to ψ_α and substituting the result into the second equation in (19) we arrive at the L which coincides exactly with that one obtained in [3].

3.2. EQUATION OF MOTION

Equations (18), (19) become very useful when the dynamics of the super-3-brane is considered. As we know from the lesson on the super-2-brane in D=4 the corresponding equation of motion becomes very simple at passage to the nonlinear realization [11, 13]. A straightforward generalization of this equation for the case of the super-3-brane in D=6 has the form

$$\nabla^\alpha\nabla_\alpha\tilde\varphi = 0\,.$$

It is a matter of the direct calculations to check that this equation of motion is really correct and can be derived from (15) with making use of the equivalence relations (18), (19) and the Lagrangian density

$$L = \varphi\bar\varphi + 1/16\,(D\varphi)^2(\bar D\bar\varphi)^2 Z\,,\ \bar Z = Z\,,$$

where

$$Z^{-1} = 1/2\,(1 + 1/2\,A + \sqrt{1 + A + B})\,,$$
$$A = -2(\partial^{\alpha\dot\alpha}\varphi)(\partial_{\alpha\dot\alpha}\bar\varphi) + 1/4\,(D^2\varphi)(\bar D^2\bar\varphi)\,,$$
$$B = ((\partial^{\alpha\dot\alpha}\varphi)(\partial_{\alpha\dot\alpha}\bar\varphi))^2 - (\partial\varphi)^2(\partial\bar\varphi)^2\,.$$

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[1] Here we consider only the transformations of 2-d supersymmetry.

4. CONCLUSION

Thus, we have shown that the nontrivial generalization of the Goldstone-Maxwell of the N=2, D=4 supersymmetry can be obtained at the expense of the central charge extension of the related superalgebra (1). The main point of this generalization is the new gauge N=2 prepotential □ and □ with nontrivial dependence on the central charge (8), the irreducibility conditions (7) and the special class of the gauge transformations (11), (12). On the base of this representation we find the natural explanation of the Nambu-Goldstone mechanism of the corresponding PBGS theory of the supersymmetric 3-brane which gives the simple form of the related superfield action (14). The next advantage of this approach is the consistent way of passing to the relevant nonlinear realization. The latter allows us to answer the very important question of how to define the covariant constraints which give the possibility to express the Lagrangian density of the superbrane considered through the basic Goldstone scalar superfield φ. In this contribution we demonstrated that the only adequate answer to this question can be obtained in the framework of our approach [12], which provides the simple algorithmic procedure of imposing the corresponding covariant constraints (18) and (19). By the way, up to now such constraints was considered only for the simplest examples [9,10,14].

Unfortunately, aside of this work we left a lot of questions that would throw light on the different very intrigue problems of this theory. First, we mean the problem of violation of the SU(2) automorphisms in our approach and the possible connection of the prepotential considered with known one [16,17]. Second, from the papers [13,14,15] we know that the action of the supersymmetric 2-brane can be equally represented in terms of the basic covariant objects of the underlying nonlinear realization. It seems appropriate to ask the same question of searching the corresponding action for the super-3-brane too. We hope to turn to these and the accompanying questions in our further publications.

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