On Low-Energy Effective Action in $\mathcal{N} = 3$ Supersymmetric Gauge Theory

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Abstract

The problem of construction of low-energy effective action in $\mathcal{N} = 3$ SYM theory is considered within the harmonic superspace (HSS) approach. The low-energy effective action is supposed to be a gauge- and scale-invariant functional in $\mathcal{N} = 3$ HSS reproducing the term F^4/ϕ^4 in components. This functional is found as a scale-invariant generalization of the F^4-term in $\mathcal{N} = 3$ supersymmetric Born-Infeld action.

1 Introduction

Extended supersymmetric field theories play the important role in modern high-energy theoretical physics due to their beautiful classical and quantum properties and close relations to string/brane theory. One of the most popular examples is $\mathcal{N} = 4$ gauge theory, quantum aspects of which attracts much attention. The symmetries of this theory are so rich ($\mathcal{N} = 4$ superconformal symmetry) that many properties of this model can be proved on the symmetry grounds only.

In this paper we study the model of $\mathcal{N} = 3$ super Yang-Mills formulated in $\mathcal{N} = 3$ harmonic superspace [1]. This model, like $\mathcal{N} = 4$ SYM, is known

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to be finite [2] and superconformally invariant [3]. Moreover, $\mathcal{N}=3$ SYM model describes the dynamics of the same multiplet of physical field as $\mathcal{N}=4$ one (see e.g. book [4] for a review). Therefore it can be considered as an alternative off-shell extension of $\mathcal{N}=4$ model while on-shell they both have equivalent dynamics. However, the structure of low-energy effective action in $\mathcal{N}=3$ SYM model was not studied as yet. As another evidence in favour of usefulness of the $\mathcal{N}=3$ HSS approach, it was recently shown [5] that $\mathcal{N}=3$ SYM theory in harmonic superspace is naturally generated from superstring theory.

It is well known [6, 7] that the leading term in the low-energy effective action of $\mathcal{N}=4$ SYM model in the sector of $\mathcal{N}=2$ vector multiplet has the form $\int d^4x \frac{F^2\bar{F}^2}{(\phi\phi)^2}$, where $F_{\alpha\beta}=\partial_{\alpha\dot{\alpha}}A^{\dot{\alpha}}_{\beta}+\partial_{\beta\dot{\alpha}}A^{\dot{\alpha}}_{\alpha}$ is the Abelian strength field and ϕ is the scalar field corresponding to $\mathcal{N}=2$ vector multiplet. Complete $\mathcal{N}=4$ supersymmetric generalization of the low-energy effective action including both vector fields and hypermultiplets was given in [8], the leading bosonic component of this action is

$$\int d^4x \frac{F^2\bar{F}^2}{(\bar{\phi}_i\phi^i)^2}, \tag{1}$$

where ϕ^i is SU(3) triplet of scalar fields. Since the models of $\mathcal{N}=3$ and $\mathcal{N}=4$ SYM are equivalent on-shell, we expect that the term (1) is also leading in the effective action of $\mathcal{N}=3$ SYM model. Therefore it is interesting to find $\mathcal{N}=3$ superfield action reproducing the expression (1) in the component expansion. One can expect that such an action corresponds to the low-energy effective action of $\mathcal{N}=3$ SYM model.

An important step in understanding the possible structure of the effective action in $\mathcal{N}=3$ gauge theory was the construction of $\mathcal{N}=3$ supersymmetric Born-Infeld action [9] where it was shown that there exists a natural description of the F^4 term (and all higher-order ones) by unconstrained $\mathcal{N}=3$ superfields in harmonic superspace. We suppose that the low-energy effective action of $\mathcal{N}=3$ gauge theory should correspond to a scale invariant generalization of the corresponding BI-theory. In this work we propose some possible form of such a functional satisfying the conditions of supersymmetric, gauge and scale invariances and reproducing the term (1) in components.

This talk is a review of our recent results published in [10].

2 $\mathcal{N}=3$ multiplets and actions in harmonic superspace

The $\mathcal{N}=3$ HSS was introduced in ref. [1] to construct an off-shell superfield formulation of $\mathcal{N}=3$ SYM model. The basics of the harmonic superspace method are exposed in book [11]. Throughout this paper we follow the conventions of recent works [9, 10].

The $\mathcal{N}=3$ HSS [3, 11] is defined as the superspace with coordinates $\{Z,u\}$, where $Z=\{x^{\alpha\dot\alpha},\theta^\alpha_i,\bar\theta^{i\dot\alpha}\}$ [1] is a set of standard $\mathcal{N}=3$ coordinates and u are the harmonics parameterizing the coset $SU(3)/U(1)\times U(1)$. We consider the harmonics u^I_i and their conjugate $\bar u^i_I$ $(I=1,2,3)$ as $SU(3)$ matrices

$$u^I_i\bar u^i_J=\delta^I_J,\quad u^I_i\bar u^j_I=\delta^j_i,\quad \varepsilon^{ijk}u^1_iu^2_ju^3_k=1. \tag{2}$$

The harmonic superspace $\{Z,u\}$ contains the so called analytic subspace with the coordinates $\{\zeta_A,u\}=\{x^{\alpha\dot\alpha}_A,\theta^\alpha_2,\theta^\alpha_3,\bar\theta^{1\dot\alpha},\bar\theta^{2\dot\alpha},u\}$ where

$$x^{\alpha\dot\alpha}_A=x^{\alpha\dot\alpha}-2i(\theta^\alpha_1\bar\theta^{1\dot\alpha}-\theta^\alpha_3\bar\theta^{3\dot\alpha})\,,\quad \theta^\alpha_I=\theta^\alpha_i\bar u^i_I\,,\quad \bar\theta^{I\dot\alpha}=\bar\theta^{\dot\alpha i}u^I_i\,. \tag{3}$$

The analytic superspace plays an important role in harmonic superspace approach since it is closed under supersymmetry and $\mathcal{N}=3$ SYM action is written in analytic coordinates.

The harmonic superspace is equipped with Grassmann covariant derivatives $D^I_\alpha,\bar D_{I\dot\alpha}$ and harmonic covariant ones D^I_J which form the $su(3)$ algebra (see [3, 11] for details). For example, the derivatives which define the analytic superfields are

$$D^1_\alpha=\frac{\partial}{\partial\theta^\alpha_1},\qquad \bar D_{3\dot\alpha}=-\frac{\partial}{\partial\bar\theta^{3\dot\alpha}}. \tag{4}$$

Acting on any analytic superfield, the derivatives (4) give zero. This is similar to the chiral superfields which are annihilated by the corresponding chiral derivatives.

There are two multiplets of analytic superfields which are important for us. The first one is the multiplet of $\mathcal{N}=3$ gauge prepotentials V^1_2, V^2_3, which are complex (mutually conjugate) analytic superfields. Physical bosonic com-

[1] We denote by small Greek symbols the SL(2,C) spinor indices, $\alpha,\dot\alpha,\ldots=1,2$; the small Latin letters are $SU(3)$ indices, $i,j,\ldots=1,2,3$.

ponent fields are contained in the prepotentials as [1]

$$\begin{aligned} V_3^2 = \;& [(\bar\theta^1\bar\theta^2)u_k^2 - (\bar\theta^2)^2 u_k^1]\phi^k + \theta_3^\alpha\bar\theta^{2\dot\alpha}A_{\alpha\dot\alpha} - i\theta_2^\alpha\theta_3^\beta(\bar\theta^2)^2 H_{\alpha\beta} \\ & +\text{spinors and auxiliary fields}, \\ V_2^1 = \;& -\widetilde{(V_3^2)} = -[(\theta_2\theta_3)\bar u_2^k - (\theta_2)^2\bar u_3^k]\bar\phi_k + \theta_2^\alpha\bar\theta^{1\dot\alpha}A_{\alpha\dot\alpha} \\ & +i(\theta_2)^2\bar\theta^{1\dot\alpha}\bar\theta^{2\dot\beta}\bar H_{\dot\alpha\dot\beta} + \text{spinors and auxiliary fields}. \end{aligned} \tag{5}$$

Here, $\phi^i, \bar\phi_i$ are complex scalar fields, $A_{\alpha\dot\alpha}$ is a vector gauge field, $H_{\alpha\beta}$, $\bar H_{\dot\alpha\dot\beta}$ are the auxiliary fields which ensure the correct structure of the gauge field sector of the theory [9]. The prepotentials V_2^1, V_3^2 are used in the formulation of $\mathcal{N}=3$ SYM model in harmonic superspace. For example, the quadratic (free) $\mathcal{N}=3$ SYM action is

$$S_2[V] = -\frac{1}{4}\mathrm{tr}\int d\zeta(^{33}_{11})du\,[V_3^2 D_3^1 V_2^1 + \frac{1}{2}(D_2^1 V_3^2 - D_3^2 V_2^1)^2]. \tag{6}$$

The integration in (6) is performed over analytic superspace ($d\zeta(^{33}_{11})du$ is the integration measure on $\mathcal{N}=3$ analytic HSS).

The component form of the action S_2 in the sector of gauge fields is [9]

$$S_2 = \int d^4x[V^2 + \bar V^2 - 2(\bar V\bar F + VF) + \frac{1}{2}(F^2 + \bar F^2)], \tag{7}$$

where

$$\begin{aligned} & V_{\alpha\beta} = \tfrac{1}{4}(H_{\alpha\beta} + F_{\alpha\beta}), \qquad \bar V_{\dot\alpha\dot\beta} = \tfrac{1}{4}(\bar H_{\dot\alpha\dot\beta} + \bar F_{\dot\alpha\dot\beta}), \\ & F^2 = F^{\alpha\beta}F_{\alpha\beta}, \quad V^2 = V^{\alpha\beta}V_{\alpha\beta}, \quad FV = F^{\alpha\beta}V_{\alpha\beta}. \end{aligned} \tag{8}$$

The auxiliary fields $V_{\alpha\beta}$, $\bar V_{\dot\alpha\dot\beta}$ can be eliminated by their algebraic classical equations of motion

$$V_{\alpha\beta} = F_{\alpha\beta}, \qquad \bar V_{\dot\alpha\dot\beta} = \bar F_{\dot\alpha\dot\beta}. \tag{9}$$

As a result, the free classical action (7) takes the form of the usual Maxwell action

$$S_2 = -\frac{1}{2}\int d^4x(F^2 + \bar F^2). \tag{10}$$

Another important $\mathcal{N}=3$ multiplet is described by $\mathcal{N}=3$ superfield strengths which are expressed through prepotentials as

$$\begin{aligned} & W_{23} = \tfrac{1}{4}\bar D_{3\dot\alpha}\bar D_3^{\dot\alpha}V_2^3\,, \qquad & \bar W^{12} = -\tfrac{1}{4}D^{1\alpha}D_\alpha^1 V_1^2, \\ & W_{12} = D_1^3 W_{23}, & \bar W^{23} = -D_1^3\bar W^{12}, \\ & W_{13} = -D_1^2 W_{23}, & \bar W^{13} = D_2^3\bar W^{12}. \end{aligned} \tag{11}$$

Here V_1^2, V_2^3 are non-analytic prepotentials which are the solutions of zero-curvature equations [9]

$$D_1^2 V_2^1 = D_2^1 V_1^2, \qquad D_2^3 V_3^2 = D_3^2 V_2^3. \tag{12}$$

The superfields (11) have the following component structure in the sector of physical bosons [12]

$$\begin{aligned} W_{23} = \;& u_i^1 \phi^i(x_{A+}) + 4i\theta_2^\alpha \theta_3^\beta V_{\alpha\beta}(x_{A+}) \\ & +\text{spinors and auxiliary fields}, \\ \bar{W}^{12} = \;& \bar{u}_3^i \bar{\phi}_i(x_{A-}) + 4i\bar{\theta}^{1\dot{\alpha}} \bar{\theta}^{2\dot{\beta}} \bar{V}_{\dot{\alpha}\dot{\beta}}(x_{A-}) \\ & +\text{spinors and auxiliary fields}, \end{aligned} \tag{13}$$

where $x_{A\pm}^{\alpha\dot{\alpha}} = x_A^{\alpha\dot{\alpha}} \pm 2i\theta_2^\alpha \bar{\theta}^{2\dot{\alpha}}$.

The strength superfields (11) are used in construction of $\mathcal{N} = 3$ supersymmetric Born-Infeld action [9]. For example, the quartic term of the $\mathcal{N} = 3$ BI action is described by the following superfield action

$$S_4 = \frac{1}{32} \int d\zeta(^{33}_{11}) du \, \frac{(\bar{W}^{12} W_{23})^2}{X^2}. \tag{14}$$

This action produces the first nontrivial term

$$\frac{1}{2} \int d^4x \frac{F^2 \bar{F}^2}{X^2} \tag{15}$$

in the Born-Infeld action.

3 Construction of the leading term in $\mathcal{N} = 3$ SYM low-energy effective action

In this Section we construct a manifestly $\mathcal{N} = 3$ supersymmetric low-energy effective action containing the term F^4/ϕ^4 in the bosonic sector.

$\mathcal{N} = 3$ SYM theory is known to be a superconformal field theory [3], like the $\mathcal{N} = 4$ SYM one. Moreover, both these models describe the dynamics of the same multiplet of physical fields and therefore are on-shell equivalent [4]. The effective action of $\mathcal{N} = 3$ SYM model should be scale invariant.

The transformations of dilatations (scale invariance) and γ_5-symmetry (R-symmetry) act on the coordinates of harmonic superspace and superfield strengths as follows

$$\delta x^m_A = a x^m_A, \quad \delta\theta^\alpha_I = \frac{1}{2}(a+ib)\theta^\alpha_I, \quad \delta\bar\theta^{I\dot\alpha} = \frac{1}{2}(a-ib)\bar\theta^{I\dot\alpha}$$
$$\delta W_{IJ} = (-a+ib)W_{IJ}, \quad \delta\bar W^{IJ} = -(a+ib)\bar W^{IJ}. \tag{16}$$

We expect that a scale and γ_5-invariant generalization of the action (14) should correspond to the low-energy effective action of $\mathcal{N}=3$ SYM model. In components such an action should reproduce the scale and γ_5-invariant generalization of (15), that is (1). Note that exactly this term is leading in the low-energy effective action of $\mathcal{N}=4$ SYM model [6, 7]. Thus we wish to construct a generalization of the action (14) which would respect the scale- and γ_5-invariances.

To pass from (15) to the scale invariant component action (1), one should replace the dimensionful constant X by the function of scalar fields $(\phi^i\bar\phi_i)$. Therefore, to obtain a scale invariant generalization of the superfield action (14) we have to replace the constant X by some superfield expression having the same dimension and containing $\phi^i\bar\phi_i$ as the lowest component. The suitable expression is

$$\bar W^{IJ}W_{IJ} = \bar W^{12}W_{12} + \bar W^{23}W_{23} + \bar W^{13}W_{13}\,. \tag{17}$$

Indeed, the component expansion of this superfield starts with the scalars (see [12] for details)

$$\bar W^{IJ}W_{IJ}|_{\theta=\bar\theta=0} = \phi^i\bar\phi_i\,. \tag{18}$$

The expression (17) cannot be naively put into the integral in (14) in place of the constant X. The reason is that the superfield $(\bar W^{IJ}W_{IJ})$ is not analytic since the superfield strengths $\bar W^{23}, \bar W^{13}, W_{12}, W_{13}$ are not analytic, while the integration in (14) goes over the analytic superspace. Therefore we have to rewrite the action (14) in full $\mathcal{N}=3$ HSS and then to insert $\bar W^{IJ}W_{IJ}$ into the integral.

The action (14) in the full $\mathcal{N}=3$ HSS is written as

$$S_4 = \frac{1}{32}\int d^4x d^{12}\theta du \frac{1}{X^2}\left[\frac{(\bar D_1)^2}{4\Box}(W_{23})^2\right]\left[\frac{(D^3)^2}{4\Box}(\bar W^{12})^2\right]. \tag{19}$$

Replacing the constant X by the superfield $\bar{W}^{IJ}W_{IJ}$ in (19), we arrive at the action

$$S_4^{scale-inv} = \alpha \int d^4x d^{12}\theta du \frac{1}{(\bar{W}^{IJ}W_{IJ})^2} \times \\ \times \left[\frac{(\bar{D}_1)^2}{4\Box}(W_{23})^2\right]\left[\frac{(D^3)^2}{4\Box}(\bar{W}^{12})^2\right], \quad (20)$$

where α is some *dimensionless* constant. This constant cannot be fixed on the symmetry grounds only. One of the possible ways of finding α is a straightforward calculation of low-energy effective action in the framework of quantum field theory. Since the action (20) includes no any dimensional constants, it is scale invariant.

Let us study the component structure of the action (20). Note that the superfield strengths entering the action contain a multiplet of physical fields as well as an infinite number of auxiliary fields. We are interested in the component structure of the action (20) in the sector of scalar and vector physical fields. For this purpose we neglect all the derivatives of scalar fields and Maxwell field strength. Such an approximation is sufficient for retrieving the term F^4/ϕ^4 while going to components. Therefore we use the following ansatz for the superfield strengths

$$\begin{array}{ll} \hat{\bar{W}}^{12} = \bar{\phi}_3 + \bar{\omega}^{12}, & \hat{W}_{12} = \phi^3 + \omega_{12}, \\ \hat{\bar{W}}^{23} = \bar{\phi}_1 + \bar{\omega}^{23}, & \hat{W}_{23} = \phi^1 + \omega_{23}, \\ \hat{\bar{W}}^{13} = -\bar{\phi}_2 + \bar{\omega}^{13}, & \hat{W}_{13} = -\phi^2 + \omega_{13}, \end{array} \quad (21)$$

where

$$\begin{array}{ll} \bar{\phi}_I = \bar{u}^i_I \bar{\phi}_i, & \phi^I = u^I_i \phi^i\,, \\ \bar{\omega}^{IJ} = 4i\bar{\theta}^{I\dot{\alpha}}\bar{\theta}^{J\dot{\beta}}\bar{V}_{\dot{\alpha}\dot{\beta}}\,, & \omega_{IJ} = 4i\theta^\alpha_I\theta^\beta_J V_{\alpha\beta}. \end{array} \quad (22)$$

The symbol "hat" indicates that we consider only scalar and vector bosonic fields and discard any auxiliary fields. In the ansatz (21) and (22) the action (20) contains only local terms

$$\hat{S}_4^{scale-inv} = \alpha \int d^4x d^{12}\theta du \frac{(\theta_1)^2(\bar{\theta}^3)^2}{(\hat{\bar{W}}^{IJ}\hat{W}_{IJ})^2}(\hat{\bar{W}}^{12}\hat{W}_{23})^2. \quad (23)$$

Performing the integration over Grassmann and harmonic variables in (23), we obtain

$$\hat{S}_4^{scale-inv} = \frac{\alpha_0}{2}\int d^4x \frac{V^2\bar{V}^2}{(\phi^i\bar{\phi}_i)^2}, \quad (24)$$

where $\alpha_0 = \frac{32}{15}\alpha$.

Now we should express the auxiliary fields $V_{\alpha\beta}, \bar{V}_{\dot{\alpha}\dot{\beta}}$ through the physical field strengths $F_{\alpha\beta}, \bar{F}_{\dot{\alpha}\dot{\beta}}$ from the action

$$\hat{S}_2 + \hat{S}_4^{scale-inv} = \int d^4x \left[V^2 + \bar{V}^2 - 2(\bar{V}\bar{F} + VF) + \frac{1}{2}(F^2 + \bar{F}^2) + \frac{\alpha_0}{2}\frac{V^2\bar{V}^2}{(\phi^i\bar{\phi}_i)^2}\right]. \quad (25)$$

The corresponding equations of motion for the auxiliary fields $V_{\alpha\beta}$, $\bar{V}_{\dot{\alpha}\dot{\beta}}$ are

$$2F_{\alpha\beta} = V_{\alpha\beta}\left[2 + \frac{\alpha_0}{(\phi^i\bar{\phi}_i)^2}\bar{V}^2\right],$$
$$2\bar{F}_{\dot{\alpha}\dot{\beta}} = \bar{V}_{\dot{\alpha}\dot{\beta}}\left[2 + \frac{\alpha_0}{(\phi^i\bar{\phi}_i)^2}V^2\right]. \quad (26)$$

Eqs. (26) define the auxiliary fields $V_{\alpha\beta}, \bar{V}_{\dot{\alpha}\dot{\beta}}$ as functions of $F_{\alpha\beta}, \bar{F}_{\dot{\alpha}\dot{\beta}}$. The solution to these equations can be represented as a series over the Maxwell field strengths:

$$V_{\alpha\beta} = F_{\alpha\beta}\left[1 - \frac{\alpha_0}{2(\phi^i\bar{\phi}_i)^2}\bar{F}^2 + O(F^3)\right],$$
$$\bar{V}_{\dot{\alpha}\dot{\beta}} = \bar{F}_{\dot{\alpha}\dot{\beta}}\left[1 - \frac{\alpha_0}{2(\phi^i\bar{\phi}_i)^2}F^2 + O(F^3)\right]. \quad (27)$$

Substituting the solutions (27) into the action (24), we find

$$S_4^{scale-inv} = \frac{1}{2}\int d^4x \left[\frac{F^2\bar{F}^2}{(\phi^i\bar{\phi}_i)^2} - \frac{1}{2}\frac{F^2\bar{F}^2}{(\phi^i\bar{\phi}_i)^4}(F^2 + \bar{F}^2) + O(F^8)\right], \quad (28)$$

where we set $\alpha_0 = 1$ for simplicity. As a result, we see that the main bosonic component in the action (20) is exactly the term F^4/ϕ^4 (1) which is the first nontrivial term in the $\mathcal{N} = 3$ SYM model. The action (28) contains also all higher-order terms starting with F^6. However, the consideration of these terms requires the special attention (the corresponding analysis is performed in [10]).

Let us finish this Section with several comments concerning the superfield action (20).

- This action contains the nonlocal operator $\Box^{-1}$. However, as is shown above, the leading low-energy term in the component action is local.

- From the very beginning there is a freedom in distributing the derivatives among different factors in the actions (19) and (20). However, the local part of the action (20) actually does not depend on the specific pattern of such a distribution.

- As follows from eq. (28), the action $S_4^{scale-inv}$ contains the term (1) in its component expansion. We observe that in this expression the scalar fields appear in a single $SU(3)$ invariant combination. An analogous result was earlier obtained in ref. [8] for the full low-energy effective action of $\mathcal{N}=4$ SYM in the $\mathcal{N}=2$ HSS approach. The advantage of $\mathcal{N}=3$ formalism is that all scalar fields from the very beginning are included into a single $\mathcal{N}=3$ multiplet, while in the $\mathcal{N}=2$ superspace language the scalar fields are distributed between vector multiplet and hypermultiplet.

- The off-shell action (20) is manifestly supersymmetric, gauge invariant and scale invariant. It also respects the invariance under the γ_5 and $SU(3)$ transformations. Therefore, it can be considered as a candidate for the low-energy effective action in $\mathcal{N}=3$ SYM model.

4 Summary

In this paper we analyzed the possible off-shell structure of low-energy effective action of $\mathcal{N}=3$ SYM model written in $\mathcal{N}=3$ harmonic superspace. This action was obtained as an $\mathcal{N}=3$ superfield generalization of the term F^4/ϕ^4 which is leading in the low-energy effective action. It is written as a functional built out of the superfield strengths in full $\mathcal{N}=3$ superspace. This functional is manifestly supersymmetric, gauge invariant, scale and γ_5-invariant and corresponds to a scale invariant generalization of 4-th order term in the $\mathcal{N}=3$ supersymmetric BI action.

In conclusion, let us point out once again that the effective action (20) was found solely by employing the symmetries of the model and the requirement that it produces the F^4/ϕ^4 term in components. This action was determined up to an arbitrary numerical coefficient. The important problem now is to

reproduce the action (20) by direct quantum field theory computations in $\mathcal{N} = 3$ HSS.

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