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EFFECTIVE ACTION IN SIX-DIMENSIONAL HYPERMULTPLET THEORY COUPLED TO THE BACKGROUND VECTOR/TENSOR SYSTEM

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We consider the six-dimensional (1,0) hypermultiplet theory coupled to background vector/tensor system in harmonic superspace. An approach to calculating the superfield effective action is developed. It is shown that the classical actions for vector/tensor system are generated as the divergent parts of the effective action.

Keywords: *extended supersymmetry, harmonic superspace, vector/tensor hierarchy, effective action.*

1 Introduction

Construction of the non-Abelian $(1,0)$ and $(2,0)$ superconformal theories in $6D$ has attracted much attention for a long time (see e.g. [1, 2]). Such models can be considered as the candidates for dual gauge theories of the interacting multiple M5-branes [3] and can be related to near-horizon AdS_7 geometries. A crucial ingredient of this construction is the problem of description of the non-Abelian tensor multiplet gauge fields [4]. A solution to this problem has been found [5] in the framework of a $(1,0)$ tensor hierarchy [6] which, besides the Yang-Mills gauge field and the two-form gauge potentials of the tensor multiplet, contains the non-propagating three- and four-forms gauge potentials. Construction of the (2,0) models can be realized on the base of coupling the (1,0) non-Abelian tensor/vector models to the superconformal hypermultiplets [7]. The hypermultiplets are described by gauge non-linear sigma models [8] with a hyper-Kähler cone target space and minimal coupling to the superconformal tensor/vector models of [5]. On this road there are many open questions as to the dynamic description on the classical level, well as issues of related to the quantization of these models and the fate of the conformal and gauge symmetries at the quantum level. The most elegant technique for the study of these issues is the technique of harmonic superspace [9] which was extended to six dimensions in [10–12].

Superfield formulation of the tensor hierarchy has been studied in the paper [13] where a set of constraints on the super-$(p+1)$-form field strengths of non-Abelian super-p-form potentials in (1,0) $D6$ superspace has been proposed. These constraints restrict the field content of the super-p-forms to the fields of the non-Abelian tensor hierarchy. The superfield formulation of the tensor hierarchy sheds light on a supersymmetric structure of the theory and can serve as a base for the various generalizations. They can be useful for searching the superfield action and for studying the (2,0) superconformal theory by superspace methods. However, the superfield Lagrangian formulation of the theory under consideration has not been constructed so far.

It is known that the massless conformal $(n,0)$ superfields in six dimensions are divided into two classes: (i) the superfields whose first component carries any spin but it is an $USp(2n)$ singlet; (ii) the 'ultrashort' analytic superfields in harmonic superspace, their first component is a Lorentz scalar but it carries $USp(2n)$ indices [11]. All these superfields satisfy some superspace constrains. In this paper we explore the simplest superfields from above both classes, corresponding to the following three types of (1,0) 6D multiplets: the hypermultiplet [12], the vector multiplet [10] and the tensor multiplet [14]. The corresponding construction in the harmonic superspace a well-known.

1.1 Non-Abelian vector/tensor system

We begin with brief review of the general non-Abelian couplings of vectors and antisymmetric p-form fields in six dimensions following [5]. The (1,0) superconformal $6D$ field theory of vector/tensor system describes a hierarchy of non-Abelian scalar, vector and tensor fields $\{\phi^I, A^r_a, Y^{ij\ r}, B^I_{ab}, C_{abc\ r}, C_{abcd\ A}\}$ and their supersymmetric partners that label by indices $r = 1, \ldots, n_V$ and $I = 1, \ldots, n_T$. The full non-Abelian field strengths of vector and two-form gauge potentials are given as

$$\mathcal{F}^r_{ab} = \partial_{[a} A^r_{b]} - f_{st}{}^r A^s_a A^t_b + h^r_I B^I_{ab}, \tag{1}$$

$$\mathcal{H}^I_{abc} = \frac{1}{2}\mathcal{D}_{[a}B_{bc]} + d^I_{rs}A^r_{[a}\partial_b A^s_{c]} - \frac{1}{3}f_{pq}{}^s d^I_{rs}A^r_{[a}A^p_b A^q_{c]}$$
$$+g^{Ir}C_{abc\ r}. \quad (2)$$

Here $f_{[st]}{}^r$ are the structure constants, $d^I_{(rs)}$ are the d-symbols, defining the Chern–Simons couplings, and h^r_I, g^{Ir} are the covariantly constant tensors, defining the general Stückelberg-type couplings among forms of different degrees. Also the existence of the non-degenerate Lorenz-type metric η_{IJ}, so that $h^r_I = \eta_{IJ}g^{Jr} \equiv g^r_I$, $b_{Irs} = 2\eta_{IJ}d^J_{rs} \equiv d_{Irs}$ is assumed. The field strengths (1), (2) are defined such that they transform covariantly under the set of non-Abelian gauge transformations

$$\delta A_m = \mathcal{D}_m\Lambda^r - h^r_I\Lambda^I_m, \quad (3)$$

$$\delta B^I_{mn} = \mathcal{D}_{[m}\Lambda^I_{n]} - 2d^I_{rs}(\Lambda^r\mathcal{F}^s_{mn} - \frac{1}{2}A^r_{[m}\delta A^s_{n]}) - g^{Ir}\Lambda_{mn\ r}.$$

1.2 Harmonic superspace description of non-Abelian vector/tensor system

A complete set of the constraints on superfield strengths of the p-form potentials have been proposed in [13] in conventional 6D, (1,0) superspace. The tensor hierarchy formulated here in terms of a generalized Bianchi identities. Our aim is to reformulate these constraints in harmonic superspace and study their consistency conditions. First of all, the SYM field strength constraints $\mathcal{F}^{ij}_{\alpha\beta}{}^r = 0$ are not deformed by tensor multiplet and therefore has a solution as for the formulation of the vector multiplet in the harmonic superspace [10, 12]:

$$\{\mathcal{D}^i_\alpha, \mathcal{D}^j_\beta\} = -2i\varepsilon^{ij}\mathcal{D}_{\alpha\beta},$$
$$[\mathcal{D}^i_\gamma, \mathcal{D}_{\alpha\beta}] = -2i\varepsilon_{\alpha\beta\gamma\delta}\mathcal{W}^{i\delta}. \quad (4)$$

In the dimension 2 component of the generalized Bianchi identities we have

$$(\gamma^a)_{(\alpha\delta}\mathcal{D}^{(j}_{\beta)}\mathcal{W}^{i)\delta\ r} - 2\varepsilon^{ij}(\gamma^b)_{\alpha\beta}\mathcal{F}^r_{ab} = \frac{1}{2}\varepsilon^{ij}\gamma^a_{\alpha\beta}\Phi^I g^r_I. \quad (5)$$

These equation is equivalent to the following set of relations in the harmonic superspace formulation

$$\mathcal{D}^-_\alpha\mathcal{W}^{+\beta\ r} = \delta^\beta_\alpha(\mathcal{Y}^{+-\ r} + \frac{1}{2}\Phi^I g^r_I) + \frac{1}{2}\mathcal{F}^\beta_\alpha{}^r, \quad (6)$$
$$\frac{1}{4}(\mathcal{D}^-_\alpha\mathcal{W}^{+\alpha\ r} - \mathcal{D}^+_\alpha\mathcal{W}^{-\alpha\ r}) = \Phi^I g^r_I,$$
$$\mathcal{D}^+_\alpha\mathcal{W}^{+\beta\ r} = \delta^\beta_\alpha\mathcal{Y}^{++\ r}.$$

In the dimension 5/2 component of the generalized Bianchi identities

$$\mathcal{D}^i_\alpha\mathcal{F}^r_{ab} + i(\gamma_{[a})_{\alpha\delta}\mathcal{D}_{b]}\mathcal{W}^{i\delta\ r} = i(\gamma_{ab})^\beta_\alpha\Psi^{iI}_\beta g^r_I \quad (7)$$

determines the transformation law for the potential of 2-forms

$$\delta B^I_{ab} = i\epsilon_i\gamma_{ab}\Psi^{iI}.$$

The self-consistency conditions

$$\{\mathcal{D}^i_\alpha, \mathcal{D}^j_\beta\}\mathcal{W}^{k\delta} = -2i\varepsilon^{ij}\mathcal{D}_{\alpha\beta}\mathcal{W}^{k\delta}$$

leads to

$$\mathcal{D}^\pm_\alpha\Phi^I = 2i\Psi^{\pm I}_\alpha, \quad \mathcal{D}^+_\alpha\mathcal{Y}^{++\ r} = 0, \quad (8)$$
$$\mathcal{D}^-_\alpha\mathcal{Y}^{++\ r} = -2i(\mathcal{D}_{\alpha\beta}\mathcal{W}^{+\beta\ r} - 2\Psi^{+I}_\alpha g^r_I). \quad (9)$$

In the dimension 3 component of the generalized Bianchi identities is

$$\mathcal{D}_{[a}\mathcal{F}_{bc]} = \mathcal{H}^I_{abc}g^r_I, \quad \mathcal{H}_{abc} = (\mathcal{H}^{(+)} + \mathcal{H}^{(-)})_{abc}. \quad (10)$$

In the spinor notations it has the form

$$\frac{1}{2}\mathcal{D}_{(\alpha\delta}\tilde{\mathcal{F}}^\delta{}_{\beta)}{}^r = \frac{1}{3}\mathcal{H}^{(-)I}_{\alpha\beta}g^r_I,$$
$$\frac{1}{2}\mathcal{D}^{(\alpha\delta}\mathcal{F}^{\beta)}_\delta{}^r = \frac{1}{3}\mathcal{H}^{(+)\alpha\beta I}g^r_I. \quad (11)$$

Besides, we obtain the equations for

$$\mathcal{D}^i_\alpha\Psi^{jI}_\beta = -\frac{1}{2}\varepsilon^{ij}\mathcal{D}_{\alpha\beta}\Phi^I - \frac{1}{12}\varepsilon^{ij}\gamma^{abc}_{\alpha\beta}\mathcal{H}^{(-)I}_{abc} \quad (12)$$
$$+i\varepsilon_{\alpha\beta\gamma\delta}\mathcal{W}^{i\gamma\ r}\mathcal{W}^{j\delta\ s}d^I_{rs}.$$

The spinor derivative of the 3-rank tensor superfield $\mathcal{H}^I_{abc}$ is

$$\mathcal{D}^i_\alpha\mathcal{H}^I_{abc} = i\gamma^{[a}_{\alpha\beta}\mathcal{W}^{i\beta\ r}\mathcal{F}^{bc]\ s}d^I_{sr} + \frac{i}{2}\mathcal{D}^{[a}(\gamma^{bc]})^\beta_\alpha\Psi^{Ii}_\beta \quad (13)$$
$$-i\gamma^{abc}_{\alpha\beta}\mathcal{W}^{i\beta\ s}\Phi^J d_{Jsr}g^{Ir}.$$

This relation also determines the transformation law of the 3-form potential

$$\delta C_{abc\ r} = -i\epsilon_i\gamma_{abc}\mathcal{W}^{i\ s}\Phi^J d_{Jsr}.$$

The corresponding degrees of freedom are not dynamic since the generalized 4-form field strength satisfies the duality conditions

$$-\frac{1}{4!}\varepsilon^{abcdef}\mathcal{H}_{abcd\ r} = (\mathcal{F}^{ef\ s}\Phi^I$$
$$+ i\mathcal{W}^{i\ s}\gamma^{ef}\Psi^I_i)d_{I\ rs} \quad (14)$$

As shown in the paper [13], all other relations among the main superfield strengths and the equations of motion can be derived from these relations. These equations of motion allow the construct a component action of the theory in the form that coincides with the action of the superconformal vector/tensor system constructed in the papers [5,15]. One can show that the natural proposal for superfield form of action is written as follows

$$S = \int d\zeta^{(-4)}du(\Phi\mathcal{D}^{++}\mathcal{Y}^{++} + D^+_\alpha\Phi\mathcal{D}^{++}\mathcal{W}^{+\alpha}). \quad (15)$$

The superfields $\mathcal{W}^{+\alpha}$, $\mathcal{Y}^{++}$, Φ have been defined above. To derive the component action from the superfield action (15) we essentially used the all the above relations from this section.

2 One-loop effective action in the hypermultiplet theory

We will consider a calculation of superfield quantum effective action in the hypermultiplet theory coupled to the external field of vector/tensor system. We will show that the vector/tensor multiplet action, which is a generalization of the action [14] for the tensor multiplet, and higher derivative vector multiplet action [12] are generated as the divergent parts of the effective action. For simplicity we assume that the background fields is Abelian.

The classical conformal invariant action for a massless covariantly analytic superfield of the hypermultiplet of canonical dimension 2 coupled to a background 6D $\mathcal{N} = (1,0)$ vector/tensor system is written as

$$S = -\int du d\zeta^{(-4)} \tilde{q}^{+} \mathcal{D}^{++} q^{+} \ , \qquad (16)$$

with $\mathcal{D}^{++} = D^{++} + gV^{++}$ the analyticity-preserving covariant derivative and V^{++} the analytic potential. We want to emphasize that the superfield V^{++} here is not one for pure vector multiplet, the superfield strengths, involving the superfield V^{++}, obey the Bianchi identities which contain the superfield Φ related to tensor multiplet. As a result the action (16) describes interaction of hypermultiplet with vector/tensor system.

The hypermultiplet effective action Γ is defined by

$$e^{i\Gamma[V^{++}]} = \int \mathcal{D}q^{+}\mathcal{D}\tilde{q}^{+} \times \exp(-i\int d\zeta^{(-4)} \tilde{q}^{+}\mathcal{D}^{++}q^{+}) \ . \qquad (17)$$

The expression (17) yields

$$\Gamma[V^{++}] = i\mathrm{Tr}\ln \mathcal{D}^{++} = -i\mathrm{Tr}\ln G^{(1,1)}. \qquad (18)$$

Here $G^{(1,1)}(\zeta_1, u_1|\zeta_2, u_2) =< \tilde{q}^{+}(\zeta_1, u_1) q^{+}(\zeta_2, u_2) >$ is the superfield Green function. This Green function is analytic with respect to both arguments and satisfies the equation

$$\mathcal{D}_1^{++} G^{(1,1)}(1|2) = \delta_A^{(3,1)}(1|2) \ . \qquad (19)$$

Here $\delta_A^{(3,1)}(1|2)$ is the appropriate covariantly analytic delta-function [9]

$$\delta_A^{(q,4-q)} = (D_2^{+})^4 \delta^{14}(z_1 - z_2)\delta^{(q,-q)}(u_1, u_2). \qquad (20)$$

The formal solution to this equation can be found analogously to four-dimensional case [16–18] and looks like

$$G_\tau^{(1,1)}(1|2) = -\frac{1}{4\,\widehat{\Box}}(\mathcal{D}_1^{+})^4(\mathcal{D}_2^{+})^4 \frac{\delta^{14}(z_1 - z_2)}{(u_1^{+}u_2^{+})^3}, \qquad (21)$$

where $1/(u_1^{+}u_2^{+})^3$ a special harmonic distribution. In Eq. (21) the $\widehat{\Box}$ is the covariantly analytic d'Alembertian ($[D_\alpha^{+}, \widehat{\Box}] = 0$) which arises when $(\mathcal{D}^{+})^4(\mathcal{D}^{--})^2$ acts on the analytical superfield and has the form

$$\widehat{\Box} = -\frac{1}{8}(\mathcal{D}^{+})^4(\mathcal{D}^{--})^2| = \qquad (22)$$

$$\mathcal{D}_a\mathcal{D}^a + \mathcal{W}^{+\alpha}\mathcal{D}_\alpha^{-} + Y^{++}\mathcal{D}^{--} - \frac{1}{4}(\mathcal{D}_\alpha^{-}\mathcal{W}^{+\alpha}) - \frac{1}{2}\Phi \ .$$

Like in four- and five-dimensional cases [18] one can obtain the useful identity

$$(\mathcal{D}_1^{+})^4(\mathcal{D}_2^{+})^4\frac{1}{(u_1^{+}u_2^{+})^3} = (\mathcal{D}_1^{+})^4\{(u_1^{+}u_2^{+})(\mathcal{D}_1^{-})^4 \qquad (23)$$

$$-(u_1^{-}u_2^{+})\Delta^{--} - 4\,\widehat{\Box}\,\frac{(u_1^{-}u_2^{+})^2}{(u_1^{+}u_2^{+})}\} \ ,$$

where

$$\Delta^{--} = i\mathcal{D}^{\alpha\beta}\mathcal{D}_\alpha^{-}\mathcal{D}_\beta^{-} + 4\mathcal{W}^{-\alpha}\mathcal{D}_\alpha^{-} - (D_\alpha^{-}\mathcal{W}^{-\alpha}) \ . \qquad (24)$$

The definition (18) of the one-loop effective action is purely formal. The actual evaluation of the effective action can be done in various ways (see e.g. [17], [18]). Further we will follow [18] and use the definition

$$\Gamma(V) = \Gamma_{y=0} + \int_0^1 dy\partial_y\Gamma(yV) = \qquad (25)$$

$$-i\mathrm{Tr}\int_0^1 dy (V^{++}G^{(1,1)}(yV)),$$

where

$$\mathrm{Tr}(V^{++}G^{(1,1)}) = \int du_1 d\zeta_1^{(-4)} V^{++}(1) \times G^{(1,1)}(1|2)|_{1=2}. \qquad (26)$$

Here $G^{(1,1)}(yV)$ means the Green function depending on the background superfield yV^{++}.

The effective action in local approximation is represented as a series in powers of the background fields and their derivatives. Further we consider the calculation of the effective action on the base of superfield proper-time technique.

It is obvious that the leading non-vanishing contribution on the diagonal $(z_1, u_1) = (z_2, u_2)$ of the two-point function

$$-\frac{1}{4}\int du_1 d\zeta^{(-4)} V_1^{++}\frac{1}{\widehat{\Box}_1} \times (\mathcal{D}_1^{+})^4(\mathcal{D}_2^{+})^4\frac{\delta^{14}(z_1 - z_2)}{(u_1^{+}u_2^{+})^3}|_{1=2}, \qquad (27)$$

arises when $\mathcal{D}_1^{--}$ from $\widehat{\Box}_1$ hits on $(u_1^{+}u_2^{+})|_{u_1=u_2}$ and in addition at least eight spinor derivatives acting on the Grassmann delta-function are required to produce a non-vanishing result, $(\mathcal{D}^{-})^4(\mathcal{D}^{+})^4\delta^8(\theta_1-\theta_2)|_{\theta_1=\theta_2} = 1$.

To avoid the divergences on the intermediate steps it is necessary to introduce a regularization. We will use a variant of dimensional regularization (so called ω -regularization) accommodative for regularization of proper-time integral (see e.g. [19]). In the framework of the proper-time technique, the ω-regularized version of inverse operator is

$$-(\frac{1}{\widehat{\Box}})_{reg} = \int_0^\infty d(is)(is\mu^2)^\omega e^{is\widehat{\Box}}, \qquad (28)$$

where ω tends to zero after renormalization and μ is an arbitrary parameter of mass dimension. Taking into account the (28) and the relations (21), (25) ones get for effective action

$$\frac{1}{4}\int du_1 d\zeta_1^{(-4)} V^{++}(1) \int_0^\infty d(is)(is\mu^2)^\omega e^{is\widehat{\Box}_1} \qquad (29)$$

$$\times(\mathcal{D}_1^+)^4(\mathcal{D}_2^+)^4 \frac{1}{(u_1^+u_2^+)^3}\delta^{14}(z_1 - z_2)|_{1=2}.$$

Here $\delta^{14}(z_1-z_2) = \delta^6(x_1-x_2)\delta^4(\theta_1^+ - \theta_2^+)\delta^4(\theta_1^- - \theta_2^-)$. Now one uses the representation of the delta function

$$\delta^{14}(z_1 - z_2) = \int \frac{d^6p}{(2\pi)^6} e^{ip_a\rho^a}\delta^8(\rho_i^\alpha),$$

where

$$\rho^a = (x_1 - x_2)^a - 2i(\theta_1^+ - \theta_2^+)\gamma^a\theta^-, \quad \rho^{\alpha i} = (\theta_1 - \theta_2)^{\alpha i},$$

and $i = \pm$. In the expression (29) we commute the exponent $\exp ip_a\rho^a$ through all the operator factors to the left and then use the coincidence limit. It yields $e^{is\widehat{\Box}_1(X)} \cdot \delta^8(\theta_1 - \theta_2)$, where $X_a = \mathcal{D}_a + ip_a$, $X_\alpha^- = \mathcal{D}_\alpha^- + 2p_{\alpha\beta}\rho^{-\beta}$. To get expansion of effective action in background fields and their derivatives we should expand $e^{is\widehat{\Box}_1(X)}$ and calculate the momentum integrals. All these integrals have the standard Gauss form.

Further we will concentrate on calculating the divergent part of effective action. In the regularization scheme under consideration the divergences mean the pole terms of the form $\frac{1}{\omega}$. Expanding the $e^{is\widehat{\Box}_1(X)}$ in the (29) and leaving only the terms relating to divergences we obtain finally the divergent part of the effective action in the form

$$\Gamma_{div} = \frac{1}{4(4\pi)^3\omega}\int dud\zeta^{(-4)}V^{++}(\mathcal{D}_\alpha^+\Phi\mathcal{W}^{+\alpha} + \Phi Y^{++})$$

$$-\frac{1}{2(8\pi)^3\omega}S_{ISZ}\ . \qquad (30)$$

Here we used the conditions (9).

The divergent part of the effective action contains two contributions. The first of them is the part of the Abelian action (15) of the vector/tensor system proposed beyond. If we consider a sum of action (15) and first term in (30), we will see that this term from (30) determines a renormalization of coupling constant in the $\mathcal{D}^{++} = D^{++} + gV^{++}$. Second contribution is the Ivanov-Smilga-Zupnik higher-derivative action of the vector multiplet [12]

$$S_{ISZ} = \frac{1}{2}\int dud\zeta^{(-4)}Y^{++}Y^{++} \qquad (31)$$

one should emphasize once more that we have considered only the divergent parts of the effective action. Of course, the effective action contains the finite part, the calculation of which is extremely interesting but is a more difficult and delicate problem. As a result we see that the classical actions (15) and (31) of the Abelian theory are generated in quantum theory as the one-loop counterterms. It is worth to pointing out that the superfield calculation of divergent part of effective action is simple enough in comparison with component calculation and demonstrates a power of superfield methods.

3 Conclusion

We have studied a problem of quantum effective action in Abelian 6D, (1,0) hypermultiplet theory coupled to background vector/tensor system. The theory under consideration is formulated in superfield form in the framework of harmonic superfield approach. Such an effective action is generated by hypermultiplet loop and depends on vector and tensor multiplet superfields. The superfield proper-time technique for evaluating the effective action is developed. We have calculated the divergent part of the effective action and showed that it defines a renormalization of coupling constant. Also, it was shown that the actions of vector and tensor multiplets are generated as the divergent parts of effective action.

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ЭФФЕКТИВНОЕ ДЕЙСТВИЕ В ТЕОРИИ ШЕСТИМЕРНОГО ГИПЕРМУЛЬТИПЛЕТА, СВЯЗАННОГО С ФОНОВОЙ ВЕКТОР/ТЕНЗОРНОЙ СИСТЕМОЙ

Рассматривается теория гипермультиплета, связанного с фоновой вектор/тензорной системой в шестимерном (1,0) гармоническом суперпространстве. Развит подход к вычислению суперполевого эффективного действия. Показано, что классические суперполевые действия вектор/тензорной системы генерируются как расходящиеся части эффективного действия.

Ключевые слова: *расширенная суперсимметрия, гармоническое суперпространство, эффективное действие.*

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