We consider $D$-dimensional supersymmetric gauge theories with 8 supercharges ($D \leq 6$, $\mathcal{N}=8$) in the framework of harmonic superspaces. The effective Abelian low-energy action for $D=5$ contains the free and Chern-Simons terms. Effective $\mathcal{N}=8$ superfield actions for $D \leq 4$ can be written in terms of the superpotentials satisfying the superfield constraints and $(6-D)$-dimensional Laplace equations. The role of alternative harmonic structures is discussed.

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1 Introduction

The concept of harmonic superspace (HS) has been introduced firstly for the off-shell description of matter, gauge and supergravity superfield theories with $D=4$, $N_4=2$ supersymmetry [1, 2]. The HS approach has been used also for a consistent description of hypermultiplets and vector multiplet with $D=6$, $N_6=1$ supersymmetry [3]. It is convenient to use the total number of supercharges $\mathcal{N}=8$ for the classification of all these models in different dimensions $D$ instead of the number of spinor representations for supercharges $N_D$. The universality of harmonic superspaces is connected with the possibility of constructing $\mathcal{N}=8$ models in $D<6$ by a dimensional reduction. The HS actions for the hypermultiplets $q^+, \omega$ and the Yang-Mills prepotential $V^{++}$ can be described by universal expressions in all dimensions $D \leq 6$. Nevertheless, $\mathcal{N}=8$ supersymmetries have some specific features for each dimension based on differences in the structure of Lorentz groups $L_D$, maximum automorphism groups $R_D$ and the set of central charges $Z_D$. In particular, the alternative HS structures have been found for the case $D=2$ [4] and $D=3$ [5]. We shall study the HS constructions of low-energy effective actions for the Abelian gauge supermultiplets in $D=1, 2, 3$ and 5 in terms of the analytic prepotential $V^{++}$ and the constrained superfield $(6-D)$-component superfield strength $W(V^{++})$. The analogous problems have been discussed earlier in the framework of the component-field formalism or the formalism with the $\mathcal{N}=4$, $D=1, 2, 3$ superfields (see e.g. [6, 7, 8, 9]).
The main result of this work is a construction in the full $\mathcal{N}=8$ superspace for the dimensions $D=1, 2, 3$ and $5$ the Coulomb effective actions which contain the product of harmonic connections and the superpotentials $f_D(W)$ satisfying the $(6-D)$-dimensional Laplace equations.

Interactions of the superfields with $8$ supercharges have universal and specific features for different dimensions, however, solutions of the most geometric and dynamic problems can be simplified in the $HS$ formalism. The $(6-D)$-dimensional Laplace equations for the low-energy superpotentials follow from the gauge invariance in $HS$. The non-renormalization theorems are connected with the search of $R_D$-invariant solutions of these equations. Analysis of the equations for superpotentials is connected with the alternative $HS$ structures using different types of harmonics in the $D\leq 3$, $\mathcal{N}=8$ theories, which reflect adequately properties of duality transformations in these models.

2 Effective actions in the full and analytic superspaces

2.1. Let us firstly consider the harmonic superspace with $D=5$, $\mathcal{N}=8$ supersymmetry. The general five-dimensional superspace has the coordinates $z=(x^m, \theta_\alpha^i)$, where $m$ and $\alpha$ are the 5-vector and 4-spinor indices of the Lorentz group $L_5=SO(4,1)$, respectively, and $i$ is the 2-spinor index of the automorphism group $R_5=SU(2)$. The antisymmetric traceless 5D $\Gamma$-matrices $(\Gamma_m)_{\alpha\gamma}$ and invariant symplectic matrix $\Omega_{\alpha\rho}$ can be constructed in terms of the Weyl matrices and $\varepsilon$-symbols of $SL(2,C)$.

It is convenient to consider the bispinor representation of the 5D coordinates and partial derivatives

$$x^{\alpha\rho} = \frac{1}{2}(\Gamma_m)^{\alpha\rho} x^m, \quad \partial_{\alpha\rho} = \frac{1}{2}(\Gamma^m)_{\alpha\rho} \partial_m.$$ (1)

The basic relations between the spinor derivatives $D^k_\alpha$ in the general $D=5, \mathcal{N}=8$ superspace have the following form:

$$\{D^k_\alpha, D^l_\gamma\} = i\varepsilon^{kl}(\partial_{\alpha\gamma} + \frac{1}{2}\Omega_{\alpha\gamma} Z),$$ (2)

where $Z$ is the real central charge.

The $R_5$-invariant projections of spinor derivatives $D^\pm_\alpha = u^\pm_i D^i_\alpha$ and coordinates of the harmonic superspace $\zeta=(x^m_A, \theta^{\alpha+}, \theta^{\alpha-})$ can be defined with the help of the $SU(2)/U(1)$ harmonics $u^+_i$ by analogy with ref.[1] ($i$ is the 2-spinor index of $SU(2)$ and $q=\pm 1$ is the $U(1)$ charge). The analytic Abelian prepotential $V^{++}(\zeta, u)$ describes the 5D vector supermultiplet, which contains the
real scalar field $\phi$, the Maxwell field $A_m$, the isodoublet of 4-spinors $\lambda_i^\alpha$ and the auxiliary isotriplet $X^{ik}$. The real scalar superfield of this theory can be written in terms of the harmonic connection with the $U(1)$ charge 

$$W = -\frac{i}{2}D^\alpha D^\alpha \int du_1 \frac{V^{++}(x,u_1)}{(u^+u_1^+)^2} = -2iD^{(2)}V^{--},$$

where $(u^+u_1^+)^{-2}$ is the harmonic distribution [2]. These superfields satisfy the following constraints:

$$D^{++}V^{--} = D^{--}V^{++}, \quad D^{\pm\pm}W = 0,$$  \hspace{0.5cm} (4)

$$D^{(+2)}_{\alpha\rho}W = 0, \quad \text{where } D^{(+2)}_{\alpha\rho} = D^\alpha D^\rho - \frac{1}{4}\Omega_{\alpha\rho}D^{+\sigma}D^{\sigma}_{\sigma}. \hspace{0.5cm} (5)$$

It is readily to construct the most general low-energy effective $U(1)$-gauge action in the full $D=5$, $\mathcal{N}=8$ harmonic superspace

$$S_5 = \int d^5x d^8\theta du \ V^{++}V^{--}[g_5^{-2} + k_5W], \hspace{0.5cm} (6)$$

where $g_5$ is the coupling constant of dimension $1/2$, and $k_5$ is the dimensionless constant of the 5D Chern-Simons interaction.

The low-energy $D=4$, $\mathcal{N}=8$ effective action conserves the $SU(2)$ automorphism group and breaks the $U_R(1)$ symmetry. The corresponding $D=4$ superpotential $f(W,\bar{W})=[F(W) + \bar{F}(\bar{W})]$ satisfies the 2D Laplace equation which has only holomorphic and anti-holomorphic solutions.

### 2.2

The analogous $D=3$, $\mathcal{N}=8$ gauge theory can be constructed in the superspace with the automorphism group $R_3 = SU_l(2)\times SU_r(2)$. Coordinates of the general superspace are $z=(x^{\alpha\beta}$, $\theta^\alpha_{ia})$. The relations between spinor derivatives are

$$\{D^k_a, D^l_b\} = i\varepsilon^{kl}\varepsilon^{ab}\partial_{\alpha\beta} + i\varepsilon^{kl}\varepsilon_{\alpha\beta}Z^{ab}, \hspace{0.5cm} (7)$$

where $\partial_{\alpha\beta}=\partial/\partial x^{\alpha\beta}$ and $Z^{ab}$ are the central charges which commute with all generators except for the generators of $SU_r(2)$.

We consider here the two-component indices ($\alpha$, $\beta$ . . .) for the space-time group $SL(2,R)$, ($i$, $k$ . . .) for the group $SU_l(2)$ and ($a$, $b$ . . .) for $SU_r(2)$, respectively. We shall use the notation $u_i^{\pm} \equiv u_i^{(\pm1,0)}$ for the harmonics of the group $SU_l(2)$ and $\nu_i^{(0,\pm1)}$ for the $SU_r(2)$ harmonics, and also the notation $D^{\pm\pm}_l$ for the $l$-harmonic derivatives and $V^{\pm\pm}_l$ for the $l$-version of the 3D harmonic gauge superfields [5]. The notation with two $U(1)$ charges will be introduced for the biharmonic superfields. The gauge-covariant $SU_r(2)$-bispinor superfield of $D=3$, $\mathcal{N}=8$ gauge theory contains the corresponding harmonic connection [5]

$$W^{ab} = -iD^{+\alpha a}D^{+\beta b}_{\alpha}V^{--}_l, \hspace{0.5cm} (8)$$
where \( D^{+b}_\alpha = u^+_i D^{ib}_\alpha \). This superfield does not depend on harmonics in the Abelian case.

The \( SL(2,R) \times SU_l(2) \) invariant Coulomb effective action can be expressed in terms of the superpotential \( f_3(W^{ab}) \)

\[
S_3 = \int d^3x d^8 \theta d\nu \; V_{l}^{i+} V_{l}^{-} f_3(W^{ab}) .
\]

The gauge invariance produces the following constraint:

\[
D^{\pm c}_\alpha D^{\pm}_{c\beta} f_3(W^{ab}) = 0.
\]

This constraint is equivalent to the three-dimensional Laplace equation

\[
\Delta_w^3 f_3(W^{ab}) = 0,
\]

where \( g_3 \) is the coupling constant of dimension \( d = -1/2 \), and \( k_3 \) is the dimensionless constant of the \( \mathcal{N}=8 \) WZNW-type interaction of the vector multiplet. This interaction is well defined for nonzero values of the central-charge modulus \( |Z|_3 = \sqrt{Z^{ab} Z_{ab}} \), when its decomposition in terms of \( \hat{W}_{ab} = W_{ab} - Z_{ab} \) is nonsingular. It should be remarked that the superfield interactions of the 3D-vector multiplets with dimensionless constants (Chern-Simons terms) have been constructed earlier for the case \( \mathcal{N}=4 \) \([10]\) and \( \mathcal{N}=6 \) \([11, 5]\).

The constraints on the superfield \( W^{ab} \) can be interpreted as an alternative \( r \)-analyticity of the following projection of this superfield in the biharmonic superspace (BHS):

\[
W^{(0,2)} = v_a^{(0,1)} v_b^{(0,1)} W^{ab} = -i D^{(2,2)} V_l^{(-2,0)}(x, u) = -i \int du D^{(-2,2)} V_l^{(2,0)}(x, u).
\]

where \( V_l^{(\pm2,0)} \equiv V_l^{\pm} \) and \( D^{(\pm2,2)} = u_i^{(\pm1,0)} u_k^{(\pm1,0)} v_a^{(0,1)} v_b^{(0,1)} D^{iaa} D^{kbb} \).

The 3D-superpotential \([12]\) can be written in the form of the integral over the harmonics \( v_a^{(0,\pm1)} \):

\[
f_3(W^{ab}) = \int dv F_3[W^{(0,2)}, v^{(0,\pm1)}],
\]

where \( F_3 \) is an arbitrary function with \( q=(0,0) \). The effective action in the full superspace \([9]\) can be transformed to the equivalent representation in the \( r \)-analytic superspace

\[
S_3 = \int d^3x D^{(0,4)} dv [W^{(0,2)}]^2 F_3[W^{(0,2)}, v^{(0,\pm1)}],
\]

\[
D^{(0,-4)} = D^{(2,-2)} D^{(-2,-2)}.
\]
The mirror symmetry connects the $l$-vector multiplet $W^{(0,2)}(V^{(2,0)})$ with the $r$-analytic hypermultiplet $\omega_r$.

2.3. The two-dimensional $(4,4)$ superfields and the corresponding $\sigma$-models have been discussed in the ordinary superspace \[12, 13\] and in the framework of alternative harmonic superspaces \[4\]. The $(4,4)$ gauge theory has been considered in the formalism of the $(2,2)$ superspace \[7\]. We shall study the geometry of this theory in the manifestly covariant harmonic formalism. In the $(4,4)$ superspace with the coordinates $(y, \bar{y}, \theta^i\alpha, \bar{\theta}^i\bar{\alpha})$, we shall use the automorphism group $R_2 = SU_c(2) \times SU_l(2) \times SU_r(2)$ (the notation of spinor indices for these groups: $c)$ $i, k, \ldots$; $l)$ $\alpha, \beta \ldots$ and $r)$ $a, b \ldots$, respectively). The algebra of spinor derivatives in this superspace is

$$\{D_k\alpha, D_l\beta\} = \varepsilon_{kl}\varepsilon_{\alpha\beta}\partial_y, \quad (16)$$

$$\{\bar{D}_k\alpha, \bar{D}_l\beta\} = \varepsilon_{kl}\varepsilon_{\alpha\beta}\bar{\partial}_y, \quad (17)$$

$$\{D_k\alpha, \bar{D}_l\beta\} = i\varepsilon_{kl}Z_{\alpha\beta}, \quad (18)$$

where $Z_{\alpha\beta}$ are the central charges.

The superfield constraints of the nonabelian $(4,4)$ gauge theory are

$$\{\nabla_k\alpha, \nabla_l\beta\} = \varepsilon_{kl}\varepsilon_{\alpha\beta}\nabla_y, \quad (19)$$

$$\{\bar{\nabla}_k\alpha, \bar{\nabla}_l\beta\} = \varepsilon_{kl}\varepsilon_{\alpha\beta}\bar{\nabla}_y, \quad (20)$$

$$\{\nabla_k\alpha, \bar{\nabla}_l\beta\} = i\varepsilon_{kl}W_{\alpha\beta}, \quad (21)$$

The authors of ref.\[4\] have discussed the three types of harmonics: $u_i^\pm = u_i^{(+1,0,0)}$ for $SU_c(2)/U_c(1)$; $l_{\alpha}^{(0,\pm 1,0)}$ for $SU_l(2)/U_l(1)$; and $r_{a}^{(0,0,\pm 1)}$ for $SU_r(2)/U_r(1)$ (in our notation). The basic geometric structures of the gauge theory are connected mainly with the harmonics $u_i^\pm$ and the corresponding analytic coordinates $\zeta_c = (y_c, \theta^{+\alpha})$ and $\bar{\zeta}_c = (\bar{y}_c, \bar{\theta}^{+\alpha})$. The $SO(4)$-vector superfield strength $w_m$ for the $2D$ analytic gauge prepotential $V_c^{++}(\zeta_c, \bar{\zeta}_c, u)$ can be constructed by analogy with $D=3$

$$W_{\alpha\beta} \equiv (\sigma^m)_{\alpha\beta}W_m = -iD^+\bar{D}^+V^{--}_c \quad (22)$$

where $(\sigma^m)_{\alpha\beta}$ are the $SO(4)$ Weyl matrices. $W_{\alpha\beta}$ satisfies the superfield constraints analogous to the constraints of the so-called twisted multiplet \[12\].

In the full $(4,4)$ superspace, one can construct the effective action of the $U(1)$ gauge theory

$$S_2 = \int d^2xd^4\theta du \, V_c^{++}V_c^{--}f_2(W_m), \quad (D^+)^2f_2(W_m) = (\bar{D}^+)^2f_2(W_m) = 0. \quad (23)$$
The general \((4,4)\) superpotential satisfies the \(4D\) Laplace equation
\[
\Delta^w_4 f_2(W_m) = 0 , \quad \Delta^w_4 = \frac{\partial}{\partial W_m} \frac{\partial}{\partial W_m} . \tag{24}
\]

The \(R_2\)-invariant solution for the \((4,4)\) superpotential is determined uniquely
\[
f^R_2(w_2) = g_2^{-2} + k_2 w_2^{-2} , \quad w_2 = \sqrt{W_m W_m} . \tag{25}
\]

The analogous function has been considered in the derivation of the \(R_2\)-invariant \((2,2)\) Kähler potential of the \(D=2, (4, 4)\) gauge theory \([7]\). The manifestly \((4,4)\) covariant formalism of the harmonic gauge theory simplifies the proof of the non-renormalization theorem.

The biharmonic representation of the \((4,4)\) superpotential is natural for the solutions of Eq.(24)
\[
f_2(W^\alpha_\alpha) = \int dldr F_2[W^{(0,1,1)}, l, r] , \quad W^{(0,1,1)} = l^{(0,1,0)}_\alpha r^{(0,0,1)}_a W^\alpha_a , \tag{26}
\]
where \(F_2\) is a real \(rl\)-analytic function with the zero \(U(1)\)-charges.

This projection of the vector multiplet \((22)\) satisfies the conditions of the \(rl\)-analyticity in the triharmonic superspace
\[
u^{(+1,0,0)}_i l^{(0,1,0)}_\alpha D^{i\alpha} W^{(0,1,1)} \equiv D^{(+1,1,0)}_i W^{(0,1,1)} = 0 , \tag{27}
\]
\[
u^{(+1,0,0)}_i r^{(0,0,1)}_a \bar{D}^{i\alpha} W^{(0,1,1)} \equiv \bar{D}^{(+1,0,1)}_i W^{(0,1,1)} = 0 \tag{28}
\]
and the harmonic conditions
\[
D^{(+2,0,0)}_c W^{(0,1,1)} = D^{(0,2,0)}_l W^{(0,1,1)} = D^{(0,0,2)}_r W^{(0,1,1)} = 0 \tag{29}
\]
which are analogous to the constraints on the \(q^{(1,1)}\) superfield of ref.\([4]\) (this notation does not indicate the \(U_c(1)\) charge). Note that the vector multiplet \(W^{(0,1,1)}\) contains the \(2D\) vector field instead of the auxiliary scalar component in the superfield \(q^{(1,1)}\).

Using Eqs.(23) and (22) one can obtain the following equivalent representation of the effective \((4,4)\) action in the \(rl\)-analytic superspace:
\[
S_2 = \int dldr dx D^{(1,−1,0)}_l D^{(−1,−1,0)}_r \bar{D}^{(1,0,−1)}_l \bar{D}^{(−1,0,−1)}_r [W^{(0,1,1)}]^{2} F_2[W^{(0,1,1)}, l, r] , \tag{30}
\]
where
\[
W^{(0,1,1)} = - i D^{(1,1,0)}_l \bar{D}^{(1,0,1)}_l V^{(−2,0,0)} = - i \int du D^{(−1,1,0)}_l \bar{D}^{(−1,0,1)}_l V^{(2,0,0)} . \tag{31}
\]

The action of the \(q^{(1,1)}\) multiplet with an analogous structure of the \((4,4)\) \(\sigma\)-model has been constructed in ref.\([4]\). This multiplet is dual to the \(rl\)-analytic multiplet \(\omega^{(±1,±1)}\).
2.4. The one-dimensional $\sigma$-models have been considered in the $\mathcal{N}=4$ superspace [14, 15]. Recently, this superspace has been used also for the proof of the non-renormalization theorem in the $\mathcal{N}=8$ gauge theory [6].

We shall consider the $D=1$, $\mathcal{N}=8$ superspace which is based on the automorphism group $R_i=SU_c(2)\times Spin(5)$ and has the coordinates $(t, \theta^\alpha_i)$, where $i$ is the 2-spinor index and $\alpha$ is the 4-spinor index of the group $Spin(5)=USp(4)$. The algebra of spinor derivatives is

$$\{D^k_\alpha, D^l_\rho\} = i\varepsilon^{kl}\Omega_{\alpha\rho}\partial_t + i\varepsilon^{kl}Z_{\alpha\rho},$$

(32)

where $Z_{\alpha\rho}$ are central charges.

Constraints of the 1D vector multiplet correspond to the integrability conditions of the $c$-analyticity. The analytic 1D coordinates $\zeta_c=(t_c, \theta^{+\alpha}_1)$ can be defined via the standard harmonics $u^+_i\equiv u^{(\pm 1,0)}_i$. The notation and algebra of the harmonized spinor derivatives $D^\pm_\alpha$ are similar for $D=1$ and $D=5$ cases. The superfield strength for the corresponding harmonic gauge connections $V^{++}_c$ is the 5-vector $W_m$ (or traceless bispinor $W_{\alpha\rho}$) with respect to $Spin(5)$

$$W_{\alpha\rho} \equiv \frac{1}{2}(\Gamma^m)_{\alpha\rho}W_m = -iD^{(+2)}_{\alpha\rho}V^{--}_c, \quad D^{(+2)}W_{\alpha\rho} = 0,$$

(33)

where the $Spin(5)$ $\Gamma$-matrices and notation (5) are used.

In the Abelian gauge group, four components of this bispinor are twisted superfields, for example, $D^1_1W_{13}=D^3_3W_{13}=0$.

The effective action has the following form in the full $D=1$ superspace:

$$S_1 = \int dt d^8\theta du \ V^{++}_c V^{--}_c \ f_1(W_m).$$

(34)

The gauge invariance of $S_1$ is equivalent to the 5D Laplace equation for the superpotential

$$D^{(+2)}f_1(W_m) = 0 \rightarrow \Delta^w_5 f_1(W_m) = 0.$$ (35)

The $R_1$-invariant $D=1$ superpotential

$$f_1^R(w_1) = g_1^{-2} + k_1 w_1^{-3}$$

(36)

is determined via the length of the 5-vector

$$w_1 = (W^{\rho\sigma}W_{\rho\sigma})^{1/2}.$$ (37)

The biharmonic construction of the general 5D superpotential (35) requires the use of harmonics $v^{(0,\pm 1,0)}_\alpha$, $v^{(0,0,\pm 1)}_\alpha$ of the group $USp(4)$ [16] and the corresponding harmonic projection of the bispinor superfield

$$f_1(W^{\alpha\rho}) = \int dvF_1[W^{(0,1,1)}, v^{(0,1,0)}_\alpha],$$

$$W^{(0,1,1)} = v^{(0,1,0)}_\alpha v^{(0,0,1)}_\rho W^{\alpha\rho},$$

(38)

(39)
where the real function $F_1$ of the superfield $W^{(0,1,1)}$ and $v$-harmonics is considered. The constraints of the superfield $W^{\alpha \rho}$ are equivalent to the conditions of the $v$-analyticity
\[ u_i^{(\pm 1,0,0)} v_{\alpha}^{(0,1,0)} D^{i \alpha} W^{(0,1,1)} = u_i^{(\pm 1,0,0)} v_{\alpha}^{(0,0,1)} D^{i \alpha} W^{(0,1,1)} = 0 \] (40)
together with the harmonic conditions in $u$- and $v$-variables
\[ D_c^{(\pm 2,0,0)} W^{(0,1,1)} = D_v^{(0,2,0)} W^{(0,1,1)} = 0 , \] (41)
\[ D_v^{(0,0,2)} W^{(0,1,1)} = D_v^{(0,1,1)} W^{(0,1,1)} = 0 . \] (42)

By analogy with $D=2$ and 3 we can analyze the equivalent form of $S_1$ in the $v$-analytic superspace and corresponding duality transformations.

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