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The effective action for $\mathcal{N}=2$, $d3$ super Yang-Mills-Chern-Simons-Matter theories

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We review the background field method for three-dimensional Yang-Mills and Chern-Simons models in $\mathcal{N}=2$ superspace. The background field method and heat kernel techniques are applied for evaluating the low-energy effective actions in $\mathcal{N}=2$ supersymmetric Yang-Mills and Chern-Simons models as well as in $\mathcal{N}=4$ and $\mathcal{N}=8$ SYM theories.

Keywords: *superspace, effective action, super Yang-Mills, super Chern-Simons models.*

1 Introduction

The effective field theory of the ten-dimensional superstring is type IIA and type IIB supergravity. This fact, together with considerable evidence that string theory is ultraviolet finite, encourages us to think of superstring theory as the ultraviolet completion of type II supergravity. The full underlying theory includes not just the massless modes of supergravity but also extended objects, specifically strings. Subsequently it was understood that the spectrum also importantly includes extended Dp-branes for any $p=0,1,2,\ldots,9$. It has long been appreciated that this field theories is also the effective massless theories for eleven-dimensional supergravity [1] compactified on S^1.

This relation between superstring theory and 10d supergravity provides a basis to conjecture the existence of a theory that similarly completes 11d supergravity in the ultraviolet. Indeed, it was long expected that fundamental supersymmetric membranes and 5-branes would play the role in this ultraviolet completion that strings play in completing ten-dimensional supergravities (see for example Refs. [2]). This idea was further stimulated by the discovery that when compactifying 11d supergravity, wrapped membranes naturally turned into the fundamental strings of type IIA superstring theory [3]. While it has not actually proved possible to quantise fundamental membranes and derive 11d supergravity from them, it was argued via duality [4] that there is a consistent UV completion of 11d supergravity and that stable M2-branes and M5-branes are an important part of this theory (for a recent review and complete list of references see e.g. [5]).

A key feature of modern string theory is the dynamics of multiple D-branes [6], which are described by the end points of open strings. This description provides a great deal of insight into the worldvolume dynamics of the branes, which is described by familiar classes of gauge theories augmented by higher derivative corrections. It should not come as a surprise that the dynamics of multiple membranes and also of multiple 5-branes is more complex than that of multiple D-branes [7]. In parallel with our limited understanding of everything else about M-theory, relatively little has been known about the degrees of freedom localised on membranes and 5-branes that 'live' on these branes and govern their low energy dynamics, in the limit that gravity is decoupled. Furthermore these theories must satisfy certain physical requirements due to the geometrical interpretation of branes in a higher dimensional spacetime. In the last few years, however, considerable progress has been made in understanding certain strongly coupled, maximally supersymmetric, conformal field theories in three and six dimensions on multiple membranes in M-theory [8], [9]. Such models are referred to as the Bagger-Lambert-Gustavsson (BLG) and Aharony-Bergman-Jafferis-Maldacena (ABJM) theories. ABJM models defined as three-dimensional $\mathcal{N}=6$ superconformal $U(N)\times U(N)$ Chern-Simons-matter theory with level $(k,-k)$. It is conjectured to describe N M2-branes located at the fixed point of the C^4/Z_k orbifold in the static gauge. It is also argued that the ABJM model is dual to M-theory on $AdS_4\times S^7/Z_k$ at large N. For $SU(2)\times SU(2)$ gauge group, the $\mathcal{N}=6$ supersymmetry is enhanced to $\mathcal{N}=8$ and the ABJM model coincides with the BLG model. All these new superconformal field theories involve a non-dynamical gauge field, described by a Chern-Simons like term in the Lagrangian, which is coupled to matter fields, parameterizing the degrees of freedom transverse to the worldvolume of the M2-branes. Supersymmetrising this action is most effectively done in superspace.

During the last few years, quantum aspects of $d3$ supersymmetric theories at perturbative level attracted a considerable attention. On the field theory side of the AdS/CFT correspondence, major attention is paid to the correlation functions of gauge invariant operators (see [10] for a review). Another important object containing much information about quantum aspects of a field model is the low-energy effective action. For instance, the low-energy effective action of $\mathcal{N}=4$, $d=4$ SYM model is perfectly matched with the effective action of a probe D3 brane moving on the $AdS_5 \times S^5$ background. It would be very interesting to observe similar matching between the effective action of an M2 brane on the $AdS_4 \times S^7$ background and the low-energy effective actions of ABJM-like models. It is known for a long time that the quantization of a membrane worldvolume theory is very challenging and one of difficulty is the nonlocality associated with the deformation of membrane without changing its volume. A quantum supermembrane theory faces a serious problem of quantum mechanical instability [11]. As a result, a single (quantum mechanical) super - membrane does not make sense and we get a multi-body problem in its nature, which can be regarded as the origin of the continuous spectrum. Therefore, from the field theory side, the action of M2-brane should go away from the infra-red fixed point to a nonperturbative Yang-Mills-Chern-Simons system. Unfortunately, our current understanding of the latter is very insignificant in comparison with the four-dimensional case. To fill this gap we need to develop the methods of quantum field theory for studying low-energy effective actions for various three-dimensional gauge theories.

Classical conformal invariance is of course generically violated by quantum corrections. Importantly, whenever the associated gauge group is compact, the "Chern-Simons level" k assumes discrete values for the path integral to remain invariant under global gauge transformations in the quantum theory. Next, it is well known (with or without supersymmetry) that the Chern-Simons level k cannot be renormalised other than by a finite 1-loop shift [12]. But here we encounter a miracle of 2+1 dimensions: the lagrangian with vanishing superpotential is exactly conformal even at the quantum level. This only leaves the possibility of corrections to the Kähler potential of the theory. However it can be argued that these are irrelevant in the infrared. More generally, one can raise the issue of finding UV finite three-dimensional Chern-Simons-matter models which might be as interesting as the famous $\mathcal{N}=4$, $d=4$ SYM model.

As it is known, the most powerful approach to study the quantum supersymmetric field theories is to make use of an unconstrained superfield formulation. Unfortunately, such a formulation for the $\mathcal{N}=6,8$ super Chern-Simons-matter theory is not known. The best what we know up to now is only $\mathcal{N}=3$ off-shell formulation on $\mathcal{N}=3, d3$ harmonic superspace [13]. In such a formulation three out of six or eight supersymmetries are realized off-shell while the other three or five are hidden and the supersymmetry algebra is closed only on shell. The corresponding superfield actions involve two hypermultiplet superfields in the bifundamental representations of the gauge groups and two Chern-Simons gauge superfields corresponding to the left and right gauge groups. The $\mathcal{N}=3$ superconformal invariance allows only a minimal gauge interaction of the hypermultiplets. Therefore, the $\mathcal{N}=3, d3$ harmonic superspace methods should be helpful for these considerations. Alternatively, one can study the effective action in the $\mathcal{N}=2$ superspace [14]. From the point of view of $\mathcal{N}=2, d3$ supersymmetry [15], the $\mathcal{N}=4,6,8$ Chern-Simons and super Yang-Mills theory describe coupling of the $\mathcal{N}=2$ vector multiplet to the hypermultiplet $\Phi, \bar{\Phi}$ in the adjoint representation as well as one or another set of matter hypermultiplets $Q, \bar{Q}$ in the bifundamental representation.

The aim of this paper is to discuss the background field method for $\mathcal{N}=2$ super Chern-Simons theories, study the effective action in terms of unconstrained $\mathcal{N}=2, d3$ superfields and calculate of the leading low-energy contributions to the effective action. Although the various classical and quantum aspects of $\mathcal{N}=2, d3$ supersymmetric theories were extensively studied, the superfield background field method, allowing to develop manifestly gauge invariant and $\mathcal{N}=2$ supersymmetric perturbation theory has not been formulated up to now. Just this problem is solved in the our papers [14].

2 Supersymmetric Yang-Mills-Chern-Simons models

The non-Abelian $\mathcal{N}=2$ supersymmetric Chern-Simons action was constructed in [15],

$$S_{\rm CS}^{\mathcal{N}=2} = \frac{ik}{8\pi}{\rm tr}\int_0^1 dt \times \times \int d^7z\, \bar{D}^\alpha (e^{-2tV} D_\alpha e^{2tV}) e^{-2tV}\partial_t e^{2tV} . \quad (1)$$

Here t is an auxiliary real parameter and k is an integer (Chern-Simons level). In the Wess-Zumino gauge the component field decomposition for V is given by

$$V = \theta^\alpha\bar{\theta}^\beta A_{\alpha\beta} + i\theta^\alpha\bar{\theta}_\alpha\phi + i\theta^2\bar{\theta}^\alpha\bar{\lambda}_\alpha - -i\bar{\theta}^2\theta^\alpha\lambda_\alpha + \theta^2\bar{\theta}^2 D . \quad (2)$$

Here $A_{\alpha\beta}$ is a gauge vector field, ϕ is a real scalar, λ_α is a complex spinor and D is a real auxiliary field. In the Abelian case the integration over the parameter t can be explicitly done,

$$S_{\rm CS}^{\mathcal{N}=2} = = \frac{k}{2\pi}\int d^3x (\frac{1}{2}\varepsilon^{mnp}A_m\partial_n A_p + i\lambda^\alpha\bar{\lambda}_\alpha - 2\phi D) . \quad (3)$$

The action (2) allows for the $\mathcal{N}=4$ supersymmetric extension with a chiral superfield Φ in the adjoint representation,

$$S_{\rm CS}^{\mathcal{N}=4}=S_{\rm CS}^{\mathcal{N}=2}-\frac{ik}{4\pi}{\rm tr}\int d^5z\,\Phi^2-\frac{ik}{4\pi}{\rm tr}\int d^5\bar{z}\,\bar{\Phi}^2\,. \quad (4)$$

The transformations of hidden $\mathcal{N}=2$ supersymmetry with complex spinor parameter ϵ^α read

$$\Delta_\epsilon V=\epsilon^\alpha\bar{\theta}_\alpha\Phi_c-\epsilon^\alpha\theta_\alpha\bar{\Phi}_c\,, \quad (5)$$

$$\delta_\epsilon\Phi_c=-i\bar{\epsilon}^\alpha\bar{\nabla}_\alpha G\,,\quad \delta_\epsilon\bar{\Phi}_c=-i\epsilon^\alpha\nabla_\alpha G\,.$$

It is well known that the sum of Chern-Simons (2) and Yang-Mills (6) actions

$$S_{\rm SYM}^{\mathcal{N}=2}=\frac{1}{g^2}{\rm tr}\int d^7z\,G^2=-\frac{1}{2g^2}{\rm tr}\int d^5z\,W^\alpha W_\alpha\,, \quad (6)$$

where g is the dimensionfull coupling constant, $[g]=1/2$ describes topologically massive gauge theory. Similarly (4), the $\mathcal{N}=8$ supersymmetric extension of (6) reads

$$S_{\rm SYM}^{\mathcal{N}=8}=\frac{1}{g^2}{\rm tr}\int d^7z\,\left(G^2-\frac{1}{2}e^{-2V}\bar{\Phi}^ie^{2V}\Phi_i\right) \quad (7)$$

$$+\frac{1}{12g^2}\left({\rm tr}\int d^5z\,\varepsilon^{ijk}\Phi_i[\Phi_j,\Phi_k]+c.c.\right).$$

Here Φ_i, $i=1,2,3$, is a triplet of chiral superfields.

As is well known, quantization procedure of gauge theories requires imposing a gauge which explicitly breaks the invariance of the effective action under classical gauge transformations. To keep track of the gauge invariance one is to employ the background field method which was originally introduced by DeWitt and developed in many subsequent papers. The central idea of the background field method is a decomposition of the gauge fields into classical background and quantum fields (background-quantum splitting) and imposing the gauge conditions only on the quantum ones. After integrating out quantum fields, the path integral results in the gauge invariant effective action depending on the background fields. However, the background-quantum splitting can be very non-trivial for some gauge theories and, hence, the formulation of the background field method in any concrete theory demands a special study [14].

3 Superfield effective action

We will be interested in the low-energy effective action which is a functional for the massless fields obtained by integrating out all massive ones in the functional integral. In gauge theories the separation between massless and massive fields appears usually through the Higgs mechanism. Physically interesting to consider minimal gauge symmetry breaking, ${\rm SU}(N)\to{\rm SU}(N-1)\times{\rm U}(1)$ because, from the point of view of D-branes, the corresponding effective action contains the potential which appears when one separates one D-brane from the stack.

As a result, we get the following representation for the one-loop effective action

$$e^{i\Gamma[V]}=e^{iS[V]}\times \quad (8)$$
$$\times\int\mathcal{D}v\mathcal{D}b\mathcal{D}c\mathcal{D}\varphi e^{i{\rm tr}\int d^7z\,v(-\Box_{\rm v}+H)v+iS_{\rm gh}+iS_{\rm NK}}\,.$$

Schematically, it can be written as

$$\Gamma=\Gamma_{\rm v}+\Gamma_{\rm gh}\,,\quad \Gamma_{\rm v}=\frac{i}{2}{\rm Tr}_{\rm v}\ln(\Box_{\rm v}-H)\,, \quad (9)$$

$$\Gamma_{\rm gh}=-\frac{3i}{2}{\rm Tr}_+\ln\Box_+\,.$$

The contribution $\Gamma_{\rm v}$ to the one-loop effective action comes from the quantum gauge superfield while $\Gamma_{\rm gh}$ is due to ghosts. Here ${\rm Tr}_{\rm v}$ and ${\rm Tr}_+$ are the functional traces of the operators acting in the spaces of real and chiral superfields, respectively. The operator $\Box_+$ is the covariant d'Alembertian operator acting in the space of covariantly chiral superfields which was introduced in [14],

$$\Box_+=\frac{1}{16}\bar{\mathcal{D}}^2\mathcal{D}^2= \quad (10)$$
$$=\mathcal{D}^m\mathcal{D}_m+G^2+\frac{i}{2}(\mathcal{D}^\alpha W_\alpha)+iW^\alpha\mathcal{D}_\alpha\,.$$

The explicit expressions for the traces of these operators can be found after one specifies the gauge group and the background gauge superfield. The corresponding contributions to the effective action are given by

$$\Gamma_{\rm v}=-\frac{1}{\pi}\sum_{I<J}^N\int d^7z\int_0^\infty\frac{ds}{s\sqrt{i\pi s}}\frac{\mathbf{W}_{IJ}^2\bar{\mathbf{W}}_{IJ}^2}{\mathbf{B}_{IJ}^3}e^{is\mathbf{G}_{IJ}^2} \quad (11)$$

$$\times\tanh\frac{s\mathbf{B}_{IJ}}{2}\sinh^2\frac{s\mathbf{B}_{IJ}}{2}\,,$$

$$\Gamma_{\rm gh}=-\frac{3}{2\pi}\sum_{I<J}^N\int d^7z\Big[\mathbf{G}_{IJ}\ln\mathbf{G}_{IJ} \quad (12)$$

$$+\frac{1}{4}\int_0^\infty\frac{ds}{\sqrt{i\pi s}}e^{is\mathbf{G}_{IJ}^2}\frac{\mathbf{W}_{IJ}^2\bar{\mathbf{W}}_{IJ}^2}{\mathbf{B}_{IJ}^2}\left(\frac{\tanh(s\mathbf{B}_{IJ}/2)}{s\mathbf{B}_{IJ}/2}-1\right)\Big].$$

The sum of the expressions (11) and (12) gives us the resulting one-loop effective action in the pure $\mathcal{N}=2$ SYM theory for the gauge group ${\rm SU}(N)$ spontaneously broken down to ${\rm U}(1)^{N-1}$. We point out that only the leading $\mathbf{G}\ln\mathbf{G}$ term in the $\mathcal{N}=2$ SYM effective action was obtained in [16] using the duality transformations while the explicit quantum computations allow us to find all higher-order F^{2n} terms encoded in the proper-time integrals (11) and (12).

Next, we examined the $\mathcal{N}=8$ SYM model with the classical action (7). For the gauge group ${\rm SU}(N)$ spontaneously broken down to ${\rm U}(1)^{N-1}$ the trace of the logarithm in $\Gamma_{\mathcal{N}=8} = \frac{i}{2}\mathrm{Tr}_v \ln(\Box_v + \bar{\Phi}^i\Phi_i)$ is computed by standard methods,

$$\Gamma^{\mathcal{N}=8}_{\rm SYM} = -\frac{1}{\pi}\sum_{I<J}^{N}\int d^7z\int_0^\infty \frac{ds}{s\sqrt{i\pi s}}\frac{\mathbf{W}^2_{IJ}\bar{\mathbf{W}}^2_{IJ}}{\mathbf{B}^3_{IJ}} \qquad (13)$$

$$\times e^{is(\mathbf{G}^2_{IJ}+\bar{\mathbf{\Phi}}^i_{IJ}\mathbf{\Phi}_{i\,IJ})}\tanh\frac{s\mathbf{B}_{IJ}}{2}\sinh^2\frac{s\mathbf{B}_{IJ}}{2}\,.$$

In the case when the gauge group SU(N) is spontaneously broken down to SU($N-1$)×U(1), the leading term in the effective action (13) is given by

$$\Gamma^{\mathcal{N}=8}_{SYM} = \frac{3(N-1)}{32\pi}\int d^7z\frac{\mathbf{W}^2\bar{\mathbf{W}}^2}{(\mathbf{G}^2+\bar{\mathbf{\Phi}}^i\mathbf{\Phi}_i)^{5/2}}+\ldots \qquad (14)$$

$$\sim \int d^3x\frac{(F^{mn}F_{mn})^2}{(f^{\underline{i}}f_{\underline{i}})^{5/2}}+\ldots\,,$$

where $f^{\underline{i}}$, $\underline{i}=1,2,\ldots,7$ are the seven scalar fields in the $\mathcal{N}=8$ SYM theory and dots stand for the higher-order terms. In [17] it was argued that the F^4 term in the $\mathcal{N}=8$ SYM effective action (14) is one-loop exact in the perturbation theory, but it receives instanton corrections.

Further we had to consider pure $\mathcal{N}=2$ Chern-Simons theory with the classical action (2). It is important to specify the background above which one computes the quantum corrections. Recall that in the SYM theory we used the constant field background. Such constraints provided us with a consistent quantum field theory as such a background was a solution of classical equations of motion. However, in the pure $\mathcal{N}=2$ Chern-Simons theory the equations of motion have only trivial solutions with vanishing gauge superfield strengths. Therefore, in quantizing the Chern-Simons theory we do not impose any constraints on the background and compute the leading terms in the derivative expansion of the effective action. In other words, there is no Coulomb branch and we need to study the effective action in the conformal branch when all fields are massless and we need to introduce an effective infrared cut-off m to avoid IR divergences. Then the effective action can be presented and evaluated by standard methods. The expression with the minimal number of superfield strengths reads

$$\Gamma_{CS} = -\frac{1}{256\pi m^5}\int d^7z(\mathcal{W}^{a\alpha}\mathcal{W}^b_\alpha\mathcal{W}^{c\beta}\mathcal{W}^d_\beta$$

$$-\frac{1}{2}\mathcal{W}^{a\alpha}\mathcal{W}^{b\beta}\mathcal{W}^c_\alpha\mathcal{W}^d_\beta)\mathbf{f}_{abcd}+O(m^{-6})\,,$$

where $\mathcal{W}^{a\alpha}\equiv(W^{a\alpha}-\bar{W}^{a\alpha})T_a$, $[T_a,T_b]=f_{abc}T_c$, and

$$\mathbf{f}_{a_1a_2a_3a_4} = f_{b_1a_1b_2}f_{b_2a_2b_3}f_{b_3a_3b_4}f_{b_4a_4b_1}\,.$$

Note that these terms do not have Abelian analogs.

The leading contribution from the Faddeev-Popov ghost superfields has the form of Yang-Mills action in the $\mathcal{N}=2$, $d=3$ superspace,

$$\Gamma_{\rm gh} = \frac{1}{8\pi^2 m}\mathrm{tr}\int d^7z\,G^2+O(m^{-2})\,. \qquad (15)$$

This demonstrates that the Yang-Mills term is generated in the effective action of pure $\mathcal{N}=2$ supersymmetric Chern-Simons theory by the ghost superfield loop. The appearance of this term in the effective action is not surprising as we break the conformal invariance and topological nature of pure $\mathcal{N}=2$ Chern-Simons theory. Clearly, this term vanishes on shell for $W_\alpha=0$.

4 Conclusion

We have reviewed the construction of the background field method for gauge field theories in the $\mathcal{N}=2$, $d=3$ superspace and demonstrated its power for calculating the low-energy effective action for $\mathcal{N}=2,4,8$, $d=3$ super Yang-Mills models and $\mathcal{N}=2$ super Chern-Simons model.

The background-field-depended operators of quadratic fluctuations, which represent the key elements of the background field formalism, are exactly found in the $\mathcal{N}=2$ super Yang-Mills and Chern-Simons models for arbitrary gauge superfield background.

The structure of one-loop effective action for these models is discussed in details. We have developed the $\mathcal{N}=2$, $d=3$ superfield heat kernel technique and applied it for calculating the low-energy effective actions in a form preserving manifest gauge invariance and $\mathcal{N}=2$ supersymmetry.

For constant gauge superfield background the heat kernel for the operator of quadratic fluctuations in the SYM theory was exactly found. This allows us to find the one-loop effective action in the $\mathcal{N}=2$, $\mathcal{N}=4$ and $\mathcal{N}=8$ SYM for such a background. However, in the pure $\mathcal{N}=2$ Chern-Simons theory only vanishing gauge superfield background is allowed as a solution of classical equations of motion. Therefore we compute the effective action in the $\mathcal{N}=2$ Chern-Simons theory only in the conformal branch when all the gauge degrees of freedom are massless. In this case we consider the arbitrary background superfield and compute the leading terms in the effective action of this model containing lowest number of gauge superfields.

We show that such off-shell effective action contains the Yang-Mills term which appears due to ghost superfield contributions.

The methods developed in [14] and reviewed in this paper can be applied to a wide class of three-dimensional extended supersymmetric gauge theories. For example, it would be interesting to study the higher loop low-energy effective action in three-dimensional $\mathcal{N}=2$ and $\mathcal{N}=4$ SYM theories with matter and in the models containing both SYM and Chern-Simons terms together. More importantly, it is tempting to study the low-energy effective action of ABJM-like models which could correspond to the effective action of the M2 brane on the $AdS_4 \times S^7$ background. The latter can provide one more non-trivial evidence of the AdS_4/CFT_3 correspondence.

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ЭФФЕКТИВНОЕ ДЕЙСТВИЕ $\mathcal{N}=2$, $d3$ СУПЕР ЯНГ-МИЛЛС-ЧЕРН-САЙМОНС ТЕОРИЙ

Мы делаем обзор метода фонового поля для трехмерных моделей Янга-Миллса и Черн-Саймонса в $\mathcal{N}=2$ суперпространстве. Метод фонового поля и методы теплового ядра применяются для вычисления низкоэнергетического эффективного действия в $\mathcal{N}=2$ суперсимметричных теориях Янга-Миллса и Черна-Саймонса, а также в $\mathcal{N}=4$ и $\mathcal{N}=8$ теориях супер Янга-Миллса.

Ключевые слова: *суперпространство, эффективное действие, суперсимметричные теории Янга-Миллса; суперсимметричные модели Черн-Саймонса .*

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