On $N = 2$ low energy effective actions

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Abstract
We propose a Wilsonian action compatible with special geometry and higher dimension $N = 2$ corrections, and show that the holomorphic contribution $F$ to the low energy effective action is independent of the infrared cutoff. We further show that for asymptotically free $SU(2)$ super Yang-Mills theories, the infrared cutoff can be tuned to cancel leading corrections to $F$. We also classify all local higher-dimensional contributions to the $N = 2$ superspace effective action that produce corrections to the Kähler potential when reduced to $N = 1$ superspace.

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1 Introduction

Recently, we computed the one-loop effective Kähler potential for $N = 2$ supersymmetric QCD [1]. For the 1-PI generating functional, we found terms that violated special geometry; we were able to interpret them as arising from higher dimension terms in $N = 2$ superspace. Here we examine the effects of these terms, and show that to quadratic order in fluctuations about the vacuum, they rescale the kinetic terms of the massive electrically charged fields without changing their masses. For gauge group $SU(2)$, this rescaling arises as a shift by a scheme dependent constant independent of the moduli.

In [1], we did not consider whether even higher dimensional terms in $N = 2$ superspace could give rise to corrections to the Kähler potential. Here we find all possible local higher dimension terms that can give such contributions.

We also found that a naive attempt to define a Wilsonian effective Kähler potential was incompatible with special geometry even for the massless sector, and hence, with $N = 2$ supersymmetry [1]. Here we show that some of the one-loop corrections to the effective Kähler potential give rise to infrared divergent terms in the component expansion of the 1-PI effective action, so we cannot ignore the problem. We propose a Wilsonian action with a field dependent cutoff (somewhat in the spirit of [3]) that is compatible with special geometry. We find that up to possible rescalings of the ultraviolet cutoff, the holomorphic function $\mathcal{F}$ [3] is not modified. We also find that for asymptotically free $SU(2)$ gauge theories, the constant shift in the coefficient of the charged field kinetic terms can be tuned to zero, preserving in detail all the results of [4].

2 Preliminaries

We begin with a brief review of $N = 2$ superspace (see [1, 5] for more details): the chiral spinor coordinates $\{\theta^a\} = \{\theta^{1a}, \theta^{2a}\}$ transform as a doublet under a rigid $SU(2)$ group unrelated to the gauge group, as do their complex conjugates $\{\bar{\theta}^a\}$, and the spinor derivatives corresponding to them. The rigid $SU(2)$ indices are raised and lowered by the antisymmetric invariant tensor $C_{ab}$, with $C_{12} = C^{12} = 1$. Super Yang-Mills theory is described by gauge-covariant spinor derivatives satisfying the constraints

$$
\begin{align*}
\{\nabla_{a\alpha}, \nabla_{b\beta}\} &= iC_{ab}C_{\alpha\beta} W, \\
\{\nabla^a_{\bar{\alpha}}, \nabla^b_{\bar{\beta}}\} &= iC^{ab}C_{\bar{\alpha}\bar{\beta}} W, \\
\{\nabla_{a\alpha}, \nabla^b_{\bar{\beta}}\} &= i\delta^b_a \nabla_{a\bar{\beta}}, \\
\{\nabla_{a\alpha}, \nabla_{b\beta}\} &= i\delta^b_a \nabla_{a\beta},
\end{align*}
$$

(2.1)
where \( W, \bar{W} \) are chiral (antichiral) scalar superfield strengths \( \nabla^a \bar{W} = \nabla_{\alphab} W = 0 \). The Bianchi identities imply the important relation

\[
\nabla^a \nabla_{ba} W = C_{ac} C_{bd} \bar{\nabla}^{d\alphab} \nabla_{\alphab} \bar{W} .
\]

(2.2)

The procedure for reducing \( N = 2 \) superfields and action to \( N = 1 \) form is well known: one defines \( N = 1 \) superfield components as

\[
\phi \equiv W|, \quad W_\alpha \equiv -\nabla^2 \bar{\theta}^\alpha W ,
\]

(2.3)

where the bar denotes setting \( \theta^2 = \bar{\theta}^2 = 0 \). One identifies \( \nabla_1 \) with the \( N = 1 \) covariant spinor derivative \( \nabla \), and rewrites the \( N = 2 \) integration measure in terms of the \( N = 1 \) measure and explicit derivatives \( Q \equiv \nabla_2 \) whose action on various fields is known: \( \int d^4x \; d^8\theta \to \int d^4x \; d^4\theta \; Q^2 \bar{Q}^2 \).

The leading term in a momentum expansion of the \( N = 2 \) superspace is the imaginary part of a \( (N = 2) \) chiral integral of a holomorphic function \( F(W) \)

\[
S_F = \text{Im} \int d^4x \; d^4\theta F(W) .
\]

(2.4)

It has the familiar \( N = 1 \) expansion

\[
S_F = \text{Im} \int d^4x \; d^2\theta \; \left[ \frac{1}{2} F_{AB}(Q^a W^A)(Q^b W^B) + F_A(Q^2 W^A) \right] |
+ \text{Im} \left[ \int d^4x \; d^2\theta \; \frac{1}{2} F_{AB}(\phi) W^{A\alphab} \bar{W}_\alphab + \int d^4x \; d^2\theta \; d^2\bar{\theta} F_A(\phi) \bar{\phi}^A \right] .
\]

(2.5)

The next term in a momentum expansion of the \( N = 2 \) superspace effective action for Yang-Mills theory is a full superspace integral of a real function \( H(W, \bar{W}) \)

\[
S_H = \int d^4x \; d^8\theta \; H(W, \bar{W}) .
\]

(2.6)

The \( N = 1 \) superspace expansion of this is

\[
S_H = \int d^4x \; d^2\theta \; d^2\bar{\theta} \left( g_{AB} \left[ -\frac{1}{2} \nabla^a \phi^A \nabla_{a\alphab} \bar{\phi}^B + i \bar{W}^{B\alphab} (\nabla^A \phi^a W^A + \Gamma^A_{\alphab} \nabla^a \phi^C W^D) \right.ight.
\]

\[
\left. - (f^A_{\alphab} \bar{W}^{B\alphab} \phi^C \nabla_{a\alphab} \bar{\phi}^D + f^B_{\alphab} W^{A\alphab} \phi^C \nabla_a \phi^D) \right.
\]

\[
\left. + (\nabla^2 \phi^B + \frac{1}{2} \Gamma^B_{\alphab} \bar{W}^{C\alphab} \bar{W}_\alphab) (\nabla^2 \bar{\phi}^A + \frac{1}{2} \Gamma^A_{\alphab} W^{E\alphab} W^E) \right] \right)
\]

\[
+ \frac{1}{4} R_{ABCD} (W^{A\alphab} \bar{W}^C_{\alphab} \bar{W}^{D\alphab} \bar{W}_\alphab) - i \mu_A \left( \frac{1}{2} \nabla^a W^A - f^A_{BC} \phi^B \bar{\phi}^C \right) ,
\]

(2.7)
where \( g, \Gamma, R \) are defined in terms of the partial derivatives of \( \mathcal{H} \)

\[
g_{AB} = \mathcal{H}_{AB}, \quad \Gamma^A_{BC} = g^{AB} \mathcal{H}_{BCD}, \quad R_{ABCD} = \mathcal{H}_{ACBD} - g_{EF} \Gamma^E_{AC} \Gamma^F_{BD},
\]

and \( \mu \) is the moment map defined by:

\[
f^C_{AB} \mathcal{H}_C W^B = \eta_A(W) + i \mu_A(W, \bar{W}),
\]

for \( \eta \) an arbitrary holomorphic function of \( W \).

### 3 Possible corrections

We now look for possible corrections to masses and kinetic terms of various fields. By supersymmetry, it is sufficient to study only the scalar fields \( \phi \), but as a consistency check, we consider all the bosonic fields, and find the effective corrections for the gauge bosons with field strengths \( f_{\alpha \beta} \) as well. To find the mass of \( \phi \), we must also find the corrections for the the vector auxiliary field \( D' \). We consistently ignore higher derivative terms: terms with more than two derivatives on \( \phi \) or terms with any derivatives on \( D' \) or \( f_{\alpha \beta} \); these terms are presumed to be artifacts of expanding the nonlocal effective action in powers of derivatives, and not to contribute to masses of physical states.

The last term in (2.7) is an \( N = 1 \) Kähler potential, and can give rise to corrections to the kinetic terms of the scalars \( \phi \), the terms linear in \( D' \), and the mass terms of the gauge bosons. Corrections to terms quadratic in \( D' \) and \( f_{\alpha \beta} \) arise from a variety of other terms in the effective action, and it is gratifying to see them add up to give a consistent result.

In \[1\], we argued that on dimensional grounds and by anomalous \( R \)-invariance, \( \mathcal{H} \) has the form

\[
\mathcal{H} = \mathcal{H}^0 + c \left( \ln \frac{\phi^2}{\Lambda^2} + g^0(\phi) \right) \left( \ln \frac{\bar{\phi}^2}{\Lambda^2} + \bar{g}^0(\bar{\phi}) \right),
\]

where \( \mathcal{H}^0(\phi, \bar{\phi}) \), the holomorphic function \( g^0(\phi) \), and its conjugate \( \bar{g}^0(\bar{\phi}) \) are all functions of dimensionless ratios of \( \phi \) and \( \bar{\phi} \). For the specific case of \( SU(2) \), gauge invariance and \( R \)-symmetry (which follows from the assumption that \( \mathcal{F} \) saturates the anomaly for \( \phi \to e^{i \alpha} \phi \)) imply that we can take \( g(\phi) = 0 \) and \( \mathcal{H}^0 = \mathcal{H}^0(t) \) where

\[
t \equiv \frac{\phi \cdot \bar{\phi}}{\sqrt{\phi^2 \bar{\phi}^2}}.
\]

\[^6\text{They may contribute to the masses of nonperturbative states such as monopoles \[7\].}\]
Before proceeding with our calculation, we make a crucial observation: whenever $[\phi, \bar{\phi}] = 0$, $t = 1$. In particular, this means that for all points in the moduli space vacua, $\langle t \rangle = 1$. We do not see any reason why similar properties should hold for higher rank groups.\footnote{The function $\mathcal{H}$ was introduced in \cite{6} for the massless (abelian) sector; for $SU(2)$, the previous arguments imply that it reduces to $c \ln(\phi^2/\Lambda^2) \ln(\bar{\phi}^2/\Lambda^2)$.}

We now turn to the actual calculation. We make a component expansion of (2.7) and expand it to second order in fluctuations about the vacuum: $\phi = a e + \varphi$, $\bar{\phi} = \bar{a} e + \bar{\varphi}$ where $e \equiv e^A T_A$ is a unit vector along the direction of $\langle \phi \rangle$: $\langle \phi_A \rangle = a e_A$. As explained above, we drop higher derivative terms. We find $D' \propto [\phi, \bar{\phi}]$; consequently, $D'$ is at least linear in $\varphi$. We also use $\nabla_{aa'} \phi = \nabla_{aa'} \varphi - ia [V_{aa'}, e]$, as well as the following expectation values of the derivatives of $t$:

$$\langle t \rangle = 1, \quad \langle t_A \rangle = 0, \quad \langle t_{AB} \rangle = \frac{1}{|a|^2} (\delta_{AB} - e_A e_B).$$

(3.3)

After a lengthy calculation, as expected, we find no corrections to the $U(1)$ coupling constant $\tau$, but, for the electrically charged $N = 2$ vector multiplets, though the mass is unchanged, we do find an overall shift in the coefficient of the kinetic term:

$$\text{Im} \left\langle \frac{F'}{a} \right\rangle \rightarrow \text{Im} \left\langle \frac{F'}{a} \right\rangle - 2 \mathcal{H}'(1).$$

(3.4)

In \cite{8}, a direct relation between the coefficient of the charged vector multiplet kinetic term and the BPS mass formula was found. Were this relation to persist, then the shift in (3.4) would pose great problems, as it would spoil duality; however, preliminary calculations \cite{7} indicate that it is rather the relation found in \cite{8} that breaks down. Of course, all problems would disappear if the numerical coefficient $\mathcal{H}'(1)$ were to vanish. As we shall see in section 5, $\mathcal{H}'(1)$ is sensitive to infrared divergences in the theory, and when the theory is asymptotically free, at least to 1-loop there exists an infrared regularization which does lead to $\mathcal{H}'(1) = 0$. In contrast to this infrared sensitivity of $\mathcal{H}$, $F$ is independent of the infrared cutoff up to rescalings of the ultraviolet scale $\Lambda$, and seems to be a physically sensible, scheme independent object.

For higher rank gauge groups with larger dimension moduli spaces, there exist $R$-invariant dimensionless ratios of moduli, and we expect a new complication to arise: the shift analogous to (3.4) may not be by a numerical (if scheme dependent) constant, but may vary over the moduli space. In that case, it seems unlikely that even the leading effects of $\mathcal{H}$ could be tuned to vanish.
4 Higher dimension terms and ambiguities

Is $\mathcal{H}$ unique? Are there ambiguities in it that we have not considered? More precisely, we want to know if there are other higher dimension terms that can contribute to the effective Kähler potential.

We first find a rather trivial ambiguity in $\mathcal{H}$ that we can resolve immediately: it seems possible to absorb a term proportional to the classical action in $\mathcal{H}$, that is, a contribution to the effective Kähler potential $\Delta K = \gamma \phi \cdot \bar{\phi}$ for a constant $\gamma$. For gauge group $SU(2)$, the differential equation that we solve to find $\mathcal{H}(t)$ from the effective Kähler potential (after subtracting out the $\mathcal{F}$ contribution is:

$$\frac{\Delta K}{\phi \cdot \bar{\phi}} = \mathcal{H}'(t) \frac{1 - t^2}{t}.$$  

(4.1)

For $\Delta K = 2\gamma \phi \cdot \bar{\phi}$, we appear to find a solution $\mathcal{H} = -\gamma \ln(1 - t^2)$; however, this is singular whenever $t = 1$, that is, in the vacuum, and therefore the coefficient of this contribution is unambiguously set to zero (this divergence has a different form than the infrared divergences that we discuss in the next section, and cannot be cancelled against them).\footnote{In \cite{1}, we set the coefficient to zero without analyzing it.} Indeed, since $\mathcal{H}$ cannot contribute to the Kähler potential for the massless fields, and $\phi \cdot \bar{\phi}$ does contribute, we know without any calculation that this term must be singular and cannot be included.

A true ambiguity comes from even higher dimension terms in $N = 2$ superspace whose $N = 1$ component expansions give rise to corrections to the effective Kähler potential. These can be classified, and are finite in number. The key observations are: (1) The expansion of the measure $\int d^4x d^8\theta \rightarrow \int d^4x d^4\theta Q^2 \bar{Q}^2$ gives a factor of $Q^2 \bar{Q}^2$ that must somehow be “absorbed”.

This can happen because (2) Eq. (2.1) implies $\{\nabla_\alpha, Q_\beta\} = iC_{\alpha\beta} \bar{W}$ and (3) Eq. (2.2) implies $Q^2 W = \nabla^2 W$. Thus, terms that can contribute to the effective Kähler potential must contain $1, Q \nabla, \nabla^2 \nabla^2$, where this represents any possible combination of these (in particular, the last includes possible $\nabla_{\alpha\dot{\alpha}} \nabla^{\alpha\dot{\alpha}}$ terms). All other terms are related to these by complex conjugation or integration by parts, or do not contribute to the effective Kähler potential. After integration by parts, the possible terms and the contributions that they give rise to are:

$$\mathcal{H} \rightarrow - \mathcal{H}_{\bar{A}B} f^{A}_{\bar{B}C} f^{C}_{\bar{D}E} \bar{\phi}^{B} \phi^{D} \phi^{E}$$

$$\mathcal{H}^{1}_{AB} \nabla^{a} W^{A} Q_{a} W^{B} \rightarrow - 2 \mathcal{H}^{1}_{AB} f^{A}_{CD} f^{B}_{EH} f^{E}_{FG} \bar{\phi}^{C} \phi^{D} \phi^{F} \phi^{G} \phi^{H}$$

8 In \cite{1}, we set the coefficient to zero without analyzing it.
\[ \mathcal{H}^2_{AB} \nabla^{\alpha a} W^A \nabla_{\alpha a} W^B \rightarrow 4 \mathcal{H}^2_{AB} f^A_{CF} f^B_{GJ} f^C_{DE} f^D_{HI} \phi^D \phi^E \phi^F \phi^G \phi^H \phi^I \phi^J \]
\[ \mathcal{H}^3_{AB} \nabla^{\alpha a} W^A \nabla_{\alpha a} \bar{W}^B \rightarrow -4 \mathcal{H}^3_{AB} f^A_{CF} f^B_{GJ} f^C_{DE} f^D_{HI} \phi^D \phi^E \phi^F \phi^G \phi^H \phi^I \phi^J \]
\[ \mathcal{H}^4_{ABC} \nabla^{\alpha a} W^A \nabla_\alpha W^B \nabla_\alpha \bar{W}^C \nabla_\alpha \bar{W}^D \rightarrow 4 \mathcal{H}^4_{ABC} f^A_{DG} f^B_{HJ} f^C_{IK} \phi^I \phi^J \phi^K \phi^L \]
\[ \mathcal{H}^5_{ABCD} \nabla_\alpha W^A \nabla_\alpha W^B \nabla_\alpha \bar{W}^C \nabla_\alpha \bar{W}^D \rightarrow 4 \mathcal{H}^5_{ABCD} f^A_{EF} f^B_{GH} f^C_{IJ} f^D_{KL} \phi^E \phi^F \phi^G \phi^H \phi^I \phi^J \phi^K \phi^L \]

(4.2)

These terms can also be written in manifestly SU(2) invariant notation; the only subtlety involves the \( \mathcal{H}^1 \) term. It arises from terms proportional to \( \nabla^2_{ab} W^A \nabla_\alpha W^B \nabla_\beta W^C \), and similar terms related by partial integration.

We have not analyzed the effects of these higher dimension terms in detail; though they may lead to modifications of the explicit 1-loop \( \mathcal{H}^0 \) found in [1], we believe that they do not change any of our qualitative conclusions. In particular, it is clear that they can only contribute to the massive sector, as they all involve commutators.

## 5 Infrared Issues

In all the previous sections, we have implicitly assumed that \( \mathcal{H} \) is well-defined. However, if we use the explicit 1-loop \( \mathcal{H} \) found in [1], or equivalently, look directly at the 1-loop effective Kähler potential, the contribution of the vector multiplet to the coefficient of the kinetic terms of the massive fields (\( \mathcal{H}'(1) \)) is divergent. Examining the explicit calculation [1, 9], we have a contribution

\[ \Delta K \propto \int d^2 k \left| Tr[\ln(k^2 + M) - \ln(k^2)] \right|, \]

where \( M \) is the matrix \( \{\phi, \bar{\phi}\} \); for SU(2), \( M \) has eigenvalues \( s, \frac{1}{2}(s \pm u) \), where \( s = \phi \cdot \bar{\phi} \) and \( u = \sqrt{\phi^2 \bar{\phi}^2} \). The last of these, \( \frac{1}{2}(s - u) \), vanishes in the vacuum at all points in moduli space. Thus we see that the problem is an infrared divergence. This leads us to try to define a Wilsonian effective action.

The naive procedure for doing this is to introduce an infrared cutoff \( \mu^2 \). However, as discussed in [1], such a cutoff is incompatible with special geometry: restricting to the abelian sector \( [\phi, \bar{\phi}] = 0 \), where \( \mathcal{H} \) can give only higher derivative contributions [1] and cannot contribute to the effective Kähler potential, we find a change \( s \ln s \rightarrow (\mu^2 + s) \ln(\mu^2 + s) \). This cannot be
written in the form $\text{Im}(\bar{\phi} F'(\phi))$ for any holomorphic function $F(\phi)$, and is thus incompatible with $N = 2$ supersymmetry.

We therefore propose a new Wilsonian effective action with a field-dependent cutoff (somewhat in the spirit of [2]). Since the mass-scale in the theory is set by $\langle s \rangle$, we introduce an infrared cutoff $\xi s = \xi \phi \cdot \bar{\phi}$ where $\xi$ is a dimensionless constant. Though the physical interpretation of this cutoff is somewhat unclear (we are cutting off our integration on the scale of the classical effective field $\phi$), to this order in the momentum expansion the definition is unambiguous, and to higher orders there is “ordering” ambiguity.

With this definition, we have calculated the one loop contribution to $H$. We subtract a contribution proportional to $s \ln(u/\Lambda^2)$, as this goes into $F$; this is $\xi$-independent aside from a $\xi$-dependent redefinition of $\Lambda$ which cancels terms proportional to the classical action $s$ as discussed at the beginning of the previous section. The remaining contribution to the effective action coming from integrating over the vector multiplet is

$$
\Delta K_V = -\frac{s}{(4\pi)^2} \left\{ (2\xi - 1) \ln(t) + \xi \ln(\xi) + (\xi + 1) \ln(\xi + 1) - \frac{1}{2}(2\xi + 1) \ln \frac{(2\xi + 1)^2 t^2 - 1}{4} \right. \\
- \left. \frac{1}{2t} \ln \frac{(2\xi + 1) t + 1}{(2\xi + 1) t - 1} \right\}.
$$

(5.2)

The $\xi$-dependent classical term that we have absorbed by rescaling $\Lambda$ is

$$
-\frac{2s}{(4\pi)^2} (1 + \xi \ln \xi - (\xi + 1) \ln(\xi + 1))
$$

(5.3)

Matter in the adjoint representation contributes

$$
\Delta K_{\text{adj}} = -\frac{2s}{(4\pi)^2} \ln(t).
$$

(5.4)

Note that in this case, the classical term

$$
\frac{2s}{(4\pi)^2} (1 + \xi \ln \xi - (\xi + 1) \ln(\xi + 1))
$$

(5.5)

has all the $\xi$-dependence. Finally, matter in the fundamental representation contributes

$$
\Delta K_{\text{fund}} = \frac{s}{(4\pi)^2} \left\{ 2\xi \ln(t) - \frac{1}{4}(4\xi + 1) \ln \frac{(4\xi + 1)^2 - 1}{(4\xi + 1)^2} \right. \\
- \left. \frac{1}{4t} \sqrt{t^2 - 1} \ln \frac{(4\xi + 1) t + \sqrt{t^2 - 1}}{(4\xi + 1) t - \sqrt{t^2 - 1}} \right\}.
$$

(5.6)

7
The classical term that has been absorbed is
\[
\left(\frac{1}{4}\right) \frac{2s}{(4\pi)^2} (1 + 4\xi \ln 4\xi - (4\xi + 1) \ln(4\xi + 1) + \ln 4) .
\] (5.7)

The three contributions (5.2, 5.4, 5.6) should be compared to the contribution from \(\mathcal{H}\), which is
\[
\Delta K = u\mathcal{H}'(t) \left(1 - t^2\right) .
\] (5.8)

To calculate the infrared regulated \(\mathcal{H}\), we equate these expressions, and find a first order differential equation for \(\mathcal{H}\) as in [1].

We can eliminate \(\mathcal{H}'(1)\) by fine tuning the parameter \(\xi\). To see when this can be done we take the contribution to the one-loop effective action coming from \(n\) matter multiplets in the adjoint representation and \(m\) matter multiplets in the fundamental representation. Expanding this function of \(t = 1 + \epsilon\) for small \(\epsilon\) and requiring that the \(\epsilon\) dependent piece vanishes gives us an equation for \(\xi\)
\[
\ln \left(\frac{1 + \xi}{\xi}\right) = 4 - 4n - \frac{2m(2\xi + 1)}{(4\xi + 1)}
\] (5.9)

This equation has a solution with \(\xi\) real and positive only if \(m < 4, n = 0\), that is, if the theory is asymptotically free; negative \(\xi\) doesn’t make physical sense for an infrared cutoff. For the conformally invariant cases \(m = 4, n = 0\) and \(m = 0, n = 1\), (5.9) is solved by \(\xi = \infty\); physically, this corresponds to completely suppressing all quantum corrections. Since that correctly gives \(\Delta F = 0\), this is expected.

**Acknowledgements**

It is a pleasure to thank Jan de Boer, Kostas Skenderis, Gordon Chalmers, Marc Grisaru, Erick Weinberg, Mike Dine, and Yuri Shirman for stimulating conversations. The work of MR and FG-R was supported in part by NSF grant No. PHY 9309888. The work of UL was supported in part by NFR grant No. F-AA/FU 04038-312 and by NorfA grant No. 94.35.049-O.

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