Gaugino Condensate and
Veneziano-Yankielowicz Effective Lagrangian

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Abstract

We study the supersymmetric pure Yang-Mills theory with semisimple Lie groups. We show that the general form of the gluino condensate is determined solely by the symmetries of the theory and it is in disagreement with the recently proposed existence of a conformal phase in SYM theory. We discuss the peculiarities of the Veneziano-Yankielowicz effective Lagrangian approach and explain how it is related to the calculation of the gluino condensate.

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1 Introduction

The exact solutions \cite{1,2,3} for the IR Wilsonian effective theory of N=1 supersymmetric QCD (SQCD) are striking in their self-consistency and seem to refuse any doubts of their validity. They brought about a new direction of model constructions in the high energy particle phenomenology based on the notion of dynamical supersymmetry breaking. The nonzero gaugino condensates of the corresponding supersymmetric Yang-Mills theory (SYM) play an important role in any such a model.

Recently it was suggested that the SYM theory contains extra, chirally symmetric, vacuum with zero gaugino condensate \cite{4}. The argument was based upon an analysis of the Veneziano-Yankielowicz effective Lagrangian for the SYM theory. If this phase indeed exists, then it has drastic consequences for the vacuum structure of SQCD. In particular, in the case of SU(N) SQCD with $N_f \leq N$ flavours of massless matter fields there will be an extra disjoint point (at the origin) in the quantum moduli space of vacua. This will imply that there is no dynamical supersymmetry breaking in a large set of models which are believed to have it.

In this letter this problem is studied from the point of view of the symmetries of the SYM theory. We derive the general form for a gaugino condensate and find that the existence of a chirally symmetric phase is not compatible with the symmetries of the theory. Also we discuss the applicability of the Veneziano-Yankielowicz Lagrangian to the description of the low energy SYM theories.

When this work was completed, a preprint by C. Csaki and H. Murayama \cite{5} appeared which reaches some of the conclusions of this letter.

2 SUSY Yang-Mills

The SYM Lagrangian describing the gluodynamics of gluons $A_\mu$ and gluinos $\lambda_\alpha$ with general compact gauge group has the form

$$\mathcal{L} = -\frac{1}{4g^2}G^a_{\mu\nu}G^a_{\mu\nu} + i\frac{g}{g_0^2}\bar{\lambda}^\dagger D^{\dot{\alpha}\dot{\beta}}\lambda^\beta + \frac{i\theta}{32\pi^2}G^a_{\mu\nu}\tilde{G}^a_{\mu\nu}. \quad (1)$$

where $G^a_{\mu\nu}$ is the gluon field strength tensor, $\tilde{G}^a_{\mu\nu}$ is the dual tensor and $D^{\dot{\alpha}\dot{\beta}}$ is the covariant derivative and all quantities are defined with respect to the adjoint representation of the gauge group. This Lagrangian may be written in terms of the gauge superfield $W_\alpha$ with
physical components \((\lambda_\alpha, A_\mu)\) as follows

\[
\mathcal{L} = \frac{1}{8\pi} \text{Im} \int d^2\theta \, \tau_0 W^\alpha W_\alpha ,
\]

where the bare gauge coupling \(\tau_0\) is defined to be \(\tau_0 = \frac{4\pi i}{\theta_0} + \frac{\theta_0}{2\pi}\). The model possesses a discrete global \(Z_{2C_2}\) symmetry\(^1\), a residual non-anomalous subgroup of the anomalous chiral \(U(1)\). One way to see these is to note that a chiral rotation of the gluino field becomes a symmetry if we combine it with a shift of the \(\theta\) parameter

\[
\theta \rightarrow \theta + 2C_2 \alpha, \quad \tau \rightarrow \tau + \frac{C_2}{\pi} \alpha
\]

where \(\alpha\) is a parameter of a chiral rotation. The physics of Yang-Mills theory is periodic in \(\theta\) with period \(2\pi\), no compensation in \(\theta\) is necessary if \(\alpha = k\pi/C_2\) where \(k\) is an integer.

We will calculate the gluino condensate starting from the definition:

\[
\langle \lambda^a \lambda^a \rangle = \frac{\int DA D\lambda \, \lambda^a \lambda^a e^{-\int \mathcal{L} d^4x}}{\int DA D\lambda \, e^{-\int \mathcal{L} d^4x}}.
\]

One may introduce an external current \(J\) and rewrite (4) as follows

\[
\langle \lambda \lambda \rangle = \frac{\delta}{\delta J} Z[J] \bigg|_{J=0} = \frac{\delta}{\delta J} \left[ \log \int DA D\lambda \, e^{-\int \mathcal{L} d^4x + \int \lambda \lambda J d^4x} \right] \bigg|_{J=0} .
\]

Using the symmetries of the theory we will examine \(Z[J]\). Let us discuss how one should interpret the generating functional \(Z[J]\). Suppose that \(Z[J]\) is regularized in such a way that the discrete symmetry \(Z_{2C_2}\) are maintained. Then the generating functional depends on the bare coupling constant \(\tau_0\), the cut-off \(M\) (the parameter of regularization) as well as the current. The discrete symmetry is realized as follows

\[
Z[\tau_0 + \frac{C_2}{\pi} \alpha, M, e^{-2i\alpha J}] = Z[\tau_0, M, J], \quad \alpha = \frac{k\pi}{C_2}, \quad k \in \mathbb{Z}
\]

where \(\alpha\) is a parameter of the chiral rotation

\[
\lambda \rightarrow e^{i\alpha} \lambda .
\]

Let us take the derivative with respect to \(\alpha\) on the left hand side of (6) and set \(\alpha\) equal to zero afterwards

\[
\frac{dZ}{d\alpha} \bigg|_{\alpha=0} = 0 \iff \frac{dZ}{d\tau_0} = \frac{2\pi i}{C_2} \int \langle \lambda \lambda \rangle_J J d^4x
\]

where \(\langle \lambda \lambda \rangle_J\) is defined by the following expression

\[
\langle \lambda \lambda \rangle_J = \frac{\int DA D\lambda \, \lambda \lambda e^{-\int \mathcal{L} d^4x + \int \lambda \lambda J d^4x}}{\int DA D\lambda \, e^{-\int \mathcal{L} d^4x + \int \lambda \lambda J d^4x}} , \quad J \neq 0 .
\]

\(^1\)\(C_2\) denotes the quadratic Casimir with \(C_2 = 1\) normalization for the fundamental and anti-fundamental representations.
The differential equation (8) cannot be solved directly for the gene rating functional because it is unknown how to formulate the initial-value problem (the Cauchy problem) for the functional integral. Let us take the functional derivative $\delta/\delta J$ of (8) at $J = 0$. This results in

$$\frac{d}{d\tau_0} \langle \lambda \lambda \rangle = \frac{2\pi i}{C_2} \langle \lambda \lambda \rangle.$$  \hspace{1cm} (10)

Using a dimensional argument, one can write down the general solution of (11)

$$\langle \lambda^{\alpha}(x)\lambda^{\alpha}(x) \rangle = c M^3 f(Mx) e^{\frac{2\pi i}{C_2} \tau_0},$$  \hspace{1cm} (11)

where $c$ is the constant of integration and $f$ is an arbitrary function of the dimensionless combination $Mx$. Supersymmetry requires the gluino condensate to be independent of the coordinate $x$, i.e. to be constant \[6\]. Thus the discrete symmetry $Z_{2C_2}$ and the supersymmetry determine the general form of the gluino condensate to be

$$\langle \lambda \lambda \rangle = c M^3 e^{-\frac{8\pi^2}{C_2 g_0}} e^{\frac{2\pi i k}{C_2}},$$  \hspace{1cm} (12)

Recall that the chiral rotation (7) and the shift in $\theta$ (3) leave the theory invariant if $\alpha = k\pi/C_2$, $k \in Z$. At $\theta = 0$ the general form of the gluino condensate is

$$\langle \lambda \lambda \rangle = c M^3 e^{-\frac{8\pi^2}{C_2 g_0}} e^{\frac{2\pi i k}{C_2}}, \quad k \in Z.$$  \hspace{1cm} (13)

The gluino condensate does not acquire an anomalous dimension under renormalization since $\langle \lambda \lambda \rangle$ is the lowest component of the superfield $W^2$ and the upper component of the same superfield contains the trace of the stress tensor. Thus the gluino condensate needs no subtractions and the renormalized expression for the gluino condensate is

$$\langle \lambda \lambda \rangle = c M^3 e^{-\frac{8\pi^2}{C_2 g_0}} e^{\frac{2\pi i k}{C_2}}, \quad k \in Z.$$  \hspace{1cm} (14)

where now $M$ is the mass scale (the point of renormalization) and $g$ is the renormalized coupling constant (the running coupling constant).

It is convenient to define the strong coupling scale $\Lambda$ at which the expression for the one-loop effective running coupling constant diverges\[4\]

$$\Lambda^{b_0} = M^{b_0} e^{-\frac{8\pi^2}{g_0}}$$  \hspace{1cm} (15)

where $b_0$ is the coefficient in the one-loop $\beta$-function ($b_0 = 3C_2$). In terms of the strong coupling scale the gluino condensate has the form

$$\langle \lambda \lambda \rangle = c \Lambda^3 e^{\frac{2\pi i k}{C_2}}, \quad k \in Z.$$  \hspace{1cm} (16)

\[2\]In any particular perturbative scheme a scheme-dependent constant can be added to the one-loop expression for the running coupling constant. This constant generates a constant rescaling of $\Lambda$ and can be ignored, since it can be absorbed by a redefinition of $M$.  

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If the constant $c$ is not zero then we can conclude that the SYM theory has $C^2$ different values of the gluino condensate. This statement is just a consequence of the global symmetries.

Alternatively we may derive the general form of the gluino condensate using the idea that the coupling constant $\tau_0$ can be treated as a spurion chiral superfield. Let us introduce the generating functional as follows

$$\tilde{Z}[\Phi, \bar{\Phi}, M] = \log \int DV \ e^{-\frac{1}{8\pi} \int d^4x \Im (\int d^2\theta \Phi WW)} ,$$

where $M$ is the parameter of the regularization and $\Phi$ is an external chiral field

$$\Phi = \phi + \theta^\alpha \Psi_\alpha + \theta \Theta F .$$

It is important to note that $\tilde{Z}$ is not a holomorphic function of $\Phi$ because in the definition of SYM Lagrangian we use the imaginary part. In terms of this generating functional the gluino condensate has the form

$$\langle \lambda \lambda \rangle = \frac{\delta}{\delta \Theta} \tilde{Z}[\Phi, \bar{\Phi}, M] \bigg|_{\Theta = 0, \bar{\Theta} = 0, \phi = \tau_0, \bar{\phi} = 0} .$$

Using the symmetries of the theory we can construct the generating functional $\tilde{Z}$. The discrete symmetry $Z_{2C^2}$ is realized in the following way

$$\tilde{Z}[\Phi + \frac{C^2}{\pi} \alpha, \bar{\Phi} + \frac{C^2}{\pi} \alpha, M] = \tilde{Z}[\Phi, \bar{\Phi}, M] , \quad \alpha = \frac{k\pi}{C^2} , \quad k \in \mathbb{Z} .$$

The supersymmetry is manifest in this superfield formulation. One can write the Taylor expansion for the generating functional in superspace

$$\tilde{Z}[\Phi, \bar{\Phi}, M] = \int d^4x \ d^2\theta G_1(\Phi) + \int d^4x \ d^2\bar{\theta} G_2(\bar{\Phi}) + \int d^4x \ d^2\theta \ d^2\bar{\theta} G_3(\Phi, \bar{\Phi}) ,$$

where $G_1, G_2$ and $G_3$ are some functions which can be fixed by the discrete symmetry and dimensional arguments

$$G_1(\Phi) = AM^3 e^{\frac{2\pi i}{C^2} \Phi} , \quad G_2(\bar{\Phi}) = BM^3 e^{-\frac{2\pi i}{C^2} \bar{\Phi}} , \quad G_3(\Phi, \bar{\Phi}) = M^2 f(\Phi - \bar{\Phi}) ,$$

where $A$ and $B$ are some constants and $f$ is an arbitrary function. Inserting (21) and (22) in expression (19) we can immediately obtain the general form of the gluino condensate

$$\langle \lambda \lambda \rangle = \frac{2\pi i}{C^2} A M^3 e^{\frac{2\pi i}{C^2} \tau_0} .$$

As we see, in both approaches the general form of the gluino condensate is fixed by only two tools, supersymmetry and discrete global symmetry.
The equation (10) has two solutions: (16) with $c \neq 0$ (or $A \neq 0$ in (23)) and the trivial $\langle \lambda \lambda \rangle = 0$ ($A = 0$). Can one further eliminate possibilities? Let us look closely at the phase which is described by the zero value of the gluino condensate. There are no mass parameters in this phase and it should be conformal or more precisely, superconformal, since we assume that supersymmetry is unbroken. To see this consider the dilation transformation $x' = e^{-t} x$, $\theta' = e^{-t/2} \theta$, $\bar{\theta}' = e^{-t/2} \bar{\theta}$. Classically it is a symmetry of the modified SYM Lagrangian in (17) but quantum mechanically not. The path integral measure in (17) is not invariant under such a dilation:

$$DV(x', \theta', \bar{\theta}') = DV(x, \theta, \bar{\theta}) \exp \left( \frac{3C^2}{4\pi} t \int d^4 x d^2 \theta W^2 + h.c. + O(1/M^4) \right).$$

(24)

where $O(1/M^4)$ refers to higher dimensional $D$-terms suppressed by powers of the cut-off. Therefore the $F$-term is exact and effectively one may consider the dilaton transformation as an imaginary shift in the spurious superfield $\Phi$: $\Phi \to \Phi + \frac{3C^2}{2\pi} t$ and correspondently the chiral transformation as a real shift in the spurious superfield $\Phi$: $\Phi \to \Phi + \frac{C}{\pi} \alpha$. If we consider the phase with zero gluino condensate then the $F$-terms will disappear in generating functional (21). This functional will be invariant under the full chiral transformation $U(1)$. Since the dilaton transformation is in the same supermultiplet with R-transformation the $D$-term which is not invariant under dilations has to be zero in this phase too. We see that phase with zero condensate has to be superconformal, if it exists.

In other words, the superfield $\Phi$ is the source for the anomaly multiplet $W^2$. The generating functional (21) is the functional for the connected Green functions of the composite superfield $W^2$. So the Legendre transform of (21) will give the effective action which satisfies the anomalous chiral Ward identities. In case when $F$-terms are absent in (21), the effective action is identically zero. The absence of anomalies contradicts with the defining properties of the pure SYM theory. There is no such kind of IR fixed point when anomalies disappear like in some SQCDs. The proposal by Kovner and Shifman disagrees with the global symmetries of the SYM theory.

At the end, we want to discuss the properties of R-currents in pure SYM and in conformal window of $SU(N_c)$ SQCD. We established that the phase with $\langle \lambda \lambda \rangle = 0$ should necessarily be superconformal. At such a superconformal fixed point the axial $R$-current $(J_{a\dot{a}} = \frac{2}{g^2} \text{Tr}(\lambda_a \bar{\lambda}_{\dot{a}}))$ must be conserved. Therefore in pure SYM theory there is no such kind of the fixed point as $R$-symmetry anomaly matching conditions will not be satisfied. In general not all solutions to a differential equation (whether this is the equation (10) or minimization equation on the effective potential considered in the next section) are physical. The physical ones should satisfy initial or boundary and all other types of conditions.

3The zero effective action for the multiplet of anomalies one can interpret only as absence of anomalies.
imposed on the model. And anomaly matching is the sieve in the case. It is interesting to note that in the conformal window of $SU(N_c)\text{ SQCD}$ the divergence of the axial current $J_{\alpha\dot{\alpha}}$ from the multiplet of anomalies becomes zero too. But unlike in SYM, there is a conserved anomaly-free linear combination (with coupling dependent coefficients) of anomalous currents. When the theory flows to the fixed RG point the anomaly-free combination becomes pure axial current $[7]$. Thus there is no new global symmetry appearing at a fixed point.

3 VY Lagrangian

We believe that in SYM theory only colorless asymptotic states exist and that a mass gap is dynamically generated. So at low energy one has to describe the theory using the new degrees of freedom. The gluons and gluinos must disappear from the low energy description. It is believed that the relevant composite degree of freedom can be naturally constructed $[8]$ in terms of the chiral superfield

$$ S = \frac{3}{32\pi^2} W^2 \equiv \frac{3}{32\pi^2} \text{Tr} W^2,$$

where the color trace above is in the adjoint representation. Using the path integral formalism one can write

$$ W(J_S, J_{\bar{S}}) = \log \int DVe^{-\int d^4x \mathcal{L} + \int d^2\theta d^4x S J_S + \int d^2\bar{\theta} d^4x J_{\bar{S}} S},$$

where $W(J_S, J_{\bar{S}})$ is the generating functional for the connected Green function of the composite operators. Here, $J_S$ and $J_{\bar{S}}$ are the chiral and antichiral currents

$$ J_S = J_1(x^+) + \theta^\alpha J_{2\alpha}(x^+) + \theta \theta J_3(x^+),$$

$$ J_{\bar{S}} = J_1^*(x^-) + \bar{\theta}^\dot{\alpha} \bar{J}_{2\dot{\alpha}}(x^-) + \bar{\theta} \bar{\theta} J_3^*(x^-),$$

where we use the standard notation $x^+ = x + i/2 \theta \bar{\theta}$ and $x^- = x - i/2 \theta \bar{\theta}$. The effective action $\Gamma(S, \bar{S})$ is defined in standard way as a Legendre transform of $W(J_S, J_{\bar{S}})$

$$ \Gamma(S, \bar{S}) = W(J_S, J_{\bar{S}}) - \int d^2\theta d^4x J_S S - \int d^2\bar{\theta} d^4x J_{\bar{S}} \bar{S}. $$

One can consider the Taylor expansion of the effective action in superspace

$$ \Gamma(S, \bar{S}) = \int d^2\theta d^2\bar{\theta} F_1(S, \bar{S}) + \int d^2\theta F_2(S) + \int d^2\bar{\theta} F_3(\bar{S})$$

and fix the functions $F_1, F_2$ and $F_3$ using the chiral Ward $[9]$ identities

$$ F_1(S, \bar{S}) = (S \bar{S})^{1/3} f \left( \frac{S^{1/3}}{(D^2 S^{1/3})^{1/2}}, \frac{\bar{S}^{1/3}}{(D^2 \bar{S}^{1/3})^{1/2}} \right), $$
\[
F_2(S) = \frac{C_2}{3} \left( S \ln \left( \frac{S}{\Lambda^3} \right) - S \right), \tag{32}
\]
\[
F_3(\bar{S}) = \frac{C_2}{3} \left( \bar{S} \ln \left( \frac{\bar{S}}{\Lambda^3} \right) - \bar{S} \right), \tag{33}
\]
where \( f(x, y) \) is an arbitrary function two variables, satisfying the reality condition
\[
f^*(x, y) = f(y, x). \tag{34}
\]
This is the most general solution of the problem. Starting from it one can immediately write the zero momentum form of the effective low energy Lagrangian (Veneziano-Yankielowicz Lagrangian) \[8\]
\[
\mathcal{L} = \frac{9}{a} \left( \frac{S}{\Lambda^3} \right)^{1/3} |D + \frac{C_2}{3} \left( S \ln \left( \frac{S}{\Lambda^3} - S \right) \right) |_F + \text{h.c.}, \tag{35}
\]
where \( a \) is a numerical parameter related to the wave-function renormalization \( Z(p^2) \) taken at zero momentum.

Let us discuss the relation between the Veneziano-Yankelowicz action and the calculation of the gluino condensate. The chiral field \( S \) can be written in the standard form
\[
S = \phi(x^+) + \theta^\alpha \psi_\alpha(x^+) + \theta \theta F(x^+) \tag{36}
\]
and be inserted in (35). The interesting part of the Lagrangian including the \( F \)-field is
\[
V(\phi, \phi^*, F, F^*) = -\frac{a}{9} (\phi \phi^*)^{-2/3} F F^* - \frac{C_2}{3} F \ln \frac{\phi}{\Lambda^3} - \frac{C_2}{3} F^* \ln \frac{\phi^*}{\Lambda^3}, \tag{37}
\]
where \( F \) is regarded as an auxiliary field. Elimination of the auxiliary field leads to the original VY scalar potential
\[
V(\phi, \phi^*) = \frac{1}{a} |\phi|^{4/3} \left( \ln^2 \frac{|\phi| C_2}{\Lambda^3 c_2} + C_2^2 (\arg \phi)^2 \right) \geq 0, \tag{38}
\]
which has the \( C_2 \) standard minima and one exotic at \( \phi = 0 \). This answer is natural, because we reproduce all solution of the differential equation (10). The differential equation (10) is just a consequence of the global symmetries of SYM. The potential (38) comes as a solution of Ward identities. As it was pointed out one has to look carefully at all solutions before identifying the true vacua of the theory. Note that all derivatives of the potential \( V(\phi, \phi^*) \) at \( \phi = \phi^* = 0 \) are not well defined. The potential \( V(\phi, \phi^*) \) was obtained as a solution of Ward identities which are the functional-differential equations. If the field \( \phi \) is regarded as a constant field then the Ward identities are just the differential equations. Since the derivatives are not well defined, the potential is not a solution of Ward identities at the origin. So the solution \( \phi = 0 \) is not compatible with Ward identities.
The non-constant function $f$ in the action \((30)-(33)\) leads to the propagation of the auxiliary field $F$. As was noted in \([9]\) the effective potential will include $\phi$ and $F$-fields in this case and it will not be bounded from below. One must eliminate the auxiliary field\([4]\). After their elimination they lead to a theory with nonlocal interactions since the auxiliary fields have nontrivial dynamical equations. But if we are interested mainly in the effective potential (the effective action with its external momenta set to zero), we neglect all higher order derivative terms.

We would like to finish with a few remarks about the effective Lagrangian approach. Historically, many of the non-perturbative properties was discovered by effective Lagrangian approach which is based on the construction of the most general possible form of the action satisfying the non-anomalous and anomalous Ward identities of the underlying theory. From this general form of the effective action one can extract, for example, the possible form of condensates. We have to remember that the notion of the effective Lagrangian is defined correctly\([5]\) when we “sit” in one of the vacua. In this sense the effective Lagrangian approach study a set of different Lagrangians, because the VY Lagrangian is not a singlevalued function of the fields and there no explicit $Z_{2C_2}$ symmetry. These features were pointed out in \([4]\) as defects in the VY Lagrangian. These problems automatically can be fixed if we consider the effective Lagrangian in one of $C_2$ vacua. The possible physical information can still be extracted from the VY effective Lagrangian, in particular some mass ratios of the bound states can be predicted \([10]\). In the end, this approach could be tested on the lattice.

4 Summary

We have derived the gaugino condensates for SYM theories starting from the constraints imposed by supersymmetry and other global symmetries. The result completely agrees with the exact result in SQCD. We have argued that the vacuum with zero gaugino condensate corresponding to the possible conformal phase of SYM theory is not compatible with the global symmetries of the theory. Therefore the exact results in SQCD are not modified and phenomenological models based on dynamical supersymmetry breaking remain viable. Finally we commented on the applicability of the Veneziano-Yankielowicz effective Lagrangian approach.

\(^4\)In the case of general function $f$, it is not possible to solve explicitly equations of motion for auxiliary fields. If lucky, one may utilize perturbation theory expanding in powers of some small parameter and higher derivative terms.

\(^5\)By this we mean that one keeps all usual properties of the QFT like a cluster decomposition, a spectral representation and etc (the Wightman QFT). The Legendre transform of the effective Lagrangian provides the generating functional for the connected Green functions of the relevant fundamental or composite fields.
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