The structure of the exact effective action and the quark confinement in MSSM QCD.

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Abstract

An expression for the exact (nonperturbative) effective action of $N=1$ supersymmetric gauge theories is proposed, supposing, that all particles except for the gauge bosons are massive. Analysis of its form shows, that instanton effects in the supersymmetric theories can lead to the quark confinement. The typical scale of confinement in MSSM QCD, calculated from the first principles, is in agreement with the experimental data. The proposed explanation is quite different from the dual Higgs mechanism.

1 Introduction

One of the most important unsolved problems of modern theoretical physics is the explanation of the quark confinement [1]. Usually the confinement is believed to be an essentially nonperturbative effect. Recently there is an considerable progress in understanding of the nonperturbative dynamics, caused by the paper [2]. In this paper the sum of all instanton corrections for the simplest case of $N=2$ supersymmetric Yang-Mills theory was found. Nevertheless, the model, considered by Seiberg and Witten is not physical. It is much more interesting to investigate $N=1$ supersymmetric theories, because indirect experimental data [3, 4] show the presence of $N=1$ supersymmetry in the standard model. That is why below we will consider this case.

From the other side, the presence of supersymmetry is very important and essential for the mechanism of confinement, proposed in this paper. This mechanism is quite different from the dual Higgs mechanism [5], which is usually used for the qualitative

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understanding of this phenomenon. Actually it is based on an assumption about the structure of effective potential beyond the perturbation theory and essentially uses the existence of auxiliary field $D$ in supersymmetric theories.

A great difference between QCD and Grand Unification scales, that is usually referred as the gauge hierarchy problem is an indirect implication, that the confinement is induced by instanton effects. Remind, that two quite different scales naturally arise in instanton calculations [3], their ratio being proportional to $\exp(-8\pi^2/e^2)$. However, expressions for the nonperturbative effective action of $N=1$ supersymmetric gauge theories, proposed in the literature [6, 8, 9] do not lead to the confinement.

Having analyzed instanton effects in this paper we propose a new (hypothetical) expression for the effective action (if all particles except for the gauge bosons are massive), which, in turn, leads to the confinement, in particular in the $SU(3)$-sector of the Minimal Supersymmetric Standard Model (MSSM QCD). The typical scale of confinement can be calculated from the first principles.

In the section 2 we discuss the structure of instanton contributions to the effective action and try to find their sum (i.e. the low energy effective action for $N=1$ supersymmetric Yang-Mills theory with matter, all particles except for the gauge bosons being massive) from some general arguments. The superpotential appears to depend on gauge and auxiliary fields in a very particular way. In the section 3 we argue, that such dependence produces the quark confinement in a sense, that asymptotic color states can not exist. The investigation shows, that there are two ”phases” in the model: a confining phase at distances, larger, than a critical one and a usual phase at small distances. The interquark potential is also found in this section. The scale of confinement is calculated and compared with the experiment in the section 4. The results are briefly discussed in the Conclusion.

## 2 Low energy effective action for MSSM QCD

MSSM QCD is a $N=1$ supersymmetric Yang-Mills theory with the gauge group $SU(3)$ and six matter supermultiplets. All particles, excluding the gauge bosons, are massive. In the present paper we will consider a similar theory with the $SU(N_c)$ gauge group and $N_f$ matter supermultiplets. It is described by the action

$$S = \frac{1}{16\pi} \text{tr} \ \text{Im} \left( \tau \int d^4x \ d^2\theta \ W^2 \right) + \frac{1}{4} \int d^4x \ d^4\theta \ \sum_{A=1}^{N_f} \left( \phi_A^+ e^{-2V} \phi_A^+ A + \tilde{\phi}_A^+ e^{2V} \phi_A^+ A \right) + S_m,$$

where the matter superfields $\phi$ and $\tilde{\phi}$ belong to the fundamental and antifundamental representations of the gauge group $SU(N_c)$, and $S_m$ is a sum of all mass terms.
Here we use the following notations:

\[
\phi(y, \theta) = \varphi(y) + \sqrt{2}\bar{\theta}(1 + \gamma_5)\psi(y) + \frac{1}{2}\bar{\theta}(1 + \gamma_5)\theta f(y),
\]

\[
\tilde{\phi}(y, \theta) = \tilde{\varphi}(y) + \sqrt{2}\bar{\theta}(1 + \gamma_5)\tilde{\psi}(y) + \frac{1}{2}\bar{\theta}(1 + \gamma_5)\tilde{\theta} f(y),
\]

\[
V(x, \theta) = -\frac{i}{2}\bar{\theta}\gamma^\mu\gamma_5\theta A_\mu(x) + i\sqrt{2}(\bar{\theta}\gamma_5 \lambda(x)) + \frac{i}{4}(\bar{\theta}\theta)^2 D(x),
\]

\[
W(y, \theta) = \frac{1}{2}(1 + \gamma_5)\left(i\sqrt{2}\lambda(y) + i\theta D(y) + \frac{1}{2}\sum_{\mu\nu}\theta F_{\mu\nu}(y) + \frac{1}{\sqrt{2}}\bar{\theta}(1 + \gamma_5)\theta \gamma^\mu D_\mu \lambda(y)\right),
\]

\[
\tau = \frac{\theta}{2\pi} + \frac{4\pi i}{e^2}, \quad y^\mu = x^\mu + \frac{i}{2}\bar{\theta}\gamma^\mu\gamma_5\theta,
\]

\[
\Sigma_{\mu\nu} = \frac{1}{2}(\gamma_\mu \gamma_\nu - \gamma_\nu \gamma_\mu), \quad D_\mu = \partial_\mu + i[A_\mu, ]
\]

where \(A_\mu\) is a gauge (gluon) field, \(\lambda\) is its spinor superpartner (gluino) and \(D\) is an auxiliary field.

Quarks are built from the fields \(\psi\) and \(\tilde{\psi}\) as follows:

\[
\Psi = \frac{1}{\sqrt{2}}\left((1 + \gamma_5)\psi + (1 - \gamma_5)\tilde{\psi}\right)
\]

The scalars \(\varphi\) and \(\bar{\varphi}\) are their superpartners (squarks), \(f\) and \(\tilde{f}\) are auxiliary fields.

Below the supersymmetry is considered to be broken. However, we are not interested in the concrete mechanism. It is only essential, that there are some "soft" terms, breaking the supersymmetry, for example, the gluino mass.

Because the confinement is a low energy phenomenon, we will try to construct the exact effective action below the thresholds for all massive particles \(^1\). It should satisfy the following requirements:

1. depend on the original (instead of composite) fields of the theory
2. agree with dynamical (perturbative and instanton) calculations (This requirement is much more restrictive, than the agreement with the transformation law of the collective coordinate measure, used in \([9]\)).

Of course, the exact effective action can be found only after carrying out the dynamical calculations and summing the series of instanton corrections. However, in this paper we present some arguments, which allow to suggest its form.

First, let us find the general structure of the effective action. We will assume, that it can be presented as a sum of an expression, invariant under the supersymmetry transformations, and (mass) terms, which can break the supersymmetry. We will be

\(^1\)As we will see below, this condition is not satisfied for \(u\) and \(d\) quarks. Nevertheless, it is not very essential, because the mass dependence of the effective action is not changed above the threshold.
interested only in holomorphic part of the supersymmetric terms. In order to find it note, that there is a relation between perturbative and instanton contributions \cite{6,10,11}. In particular, the requirement of "perturbative" renorminvariance of instanton contributions allows to construct exact $\beta$-functions of the supersymmetric theories.

Let us express this relation mathematically. Chose a scale $M$ and denote the value of the coupling constant at this scale by $\epsilon$. Then the one-loop result will be proportional to $-1/4\epsilon^2$, while instanton contributions will be proportional to $\exp(-8\pi^2 n/\epsilon^2)$, where $n$ is a module of a topological number. (One-loop contribution and instanton corrections are renorminvariant separately.)

The holomorphic part of the perturbative Wilson effective action can be written as

$$L_a = \frac{1}{16\pi} \text{Im} \text{tr} \int d^2 \theta W^2 \left( \frac{4\pi i}{e^2_{\text{eff}}} + \frac{\vartheta_{\text{eff}}}{2\pi} \right),$$

where $e_{\text{eff}}$ and $\vartheta_{\text{eff}}$ are renorminvariant functions of fields, the perturbative effective coupling constant and vacuum angle respectively. Denoting

$$z \equiv \exp \left[ 2\pi i \left( \frac{4\pi i}{e^2_{\text{eff}}} + \frac{\vartheta_{\text{eff}}}{2\pi} \right) \right],$$

the exact effective Lagrangian (including instanton contributions) can be presented in the following form:

$$L_a = \frac{1}{32\pi^2} \text{Im} \text{tr} \int d^2 \theta W^2 g(z) = \frac{1}{32\pi^2} \text{Im} \text{tr} \int d^2 \theta W^2 \left( -i \ln z + \sum_{n=1}^{\infty} c_n z^n \right).$$

(Let us note, that we did not use so far the constant field approximation, that is usually assumed for the derivation of exact results.)

In order to find the function $g(z)$, we will use the theorem, that the conditions

$$g(z) = -i \ln z + \sum_{k=0}^{\infty} c_n z^n, \quad \text{Im} g(z) > 0; \quad c_n \in \text{Im},$$

uniquely define its form up to a constant. (The inequality $\text{Im} g(z) > 0$ is a requirement of the positiveness of the effective charge and the condition $c_n \in \text{Im}$ follows from the structure of instanton corrections.) The constant is chosen so that the effective charge can have arbitrary real values (and, therefore, there is a point $z$, such that $g(z) = 0$). As it was shown in \cite{12}, these conditions uniquely lead to the following $z$-dependence

$$g(z) = 2\pi \tau(z^{-1/4}),$$

where the function $\tau(a)$ is Seiberg-Witten solution.
\[
\tau(a) = \left. \frac{da_D(u)}{da} \right|_{u = u(a)},
\]

the functions \(a\) and \(a_D\) being

\[
a(u) = \frac{\sqrt{2}}{\pi} \int_{-1}^{1} dx \frac{\sqrt{x - u}}{\sqrt{x^2 - 1}}, \quad a_D(u) = \frac{\sqrt{2}}{\pi} \int_{1}^{u} dx \frac{\sqrt{x - u}}{\sqrt{x^2 - 1}}.
\]

Therefore, finally we obtain

\[
L_a = \frac{1}{16\pi} \text{Im} \int d^2 \theta W^2 \tau(z^{-1/4}),
\]

where the parameter \(z\) can be approximately found in the one-loop approximation and exactly by the investigation of instanton measure similar to \([10, 11]\).

Because in the one-loop approximation the \(\beta\)-function is written as

\[
\beta(e) = -\frac{e^3}{16\pi^2} (3N_c - N_f),
\]

the parameter \(z\) should be proportional to \(M^{3N_c - N_f}\), where \(M\) is an UV-cutoff. The same result can be obtained from the expression for the instanton measure \([13]\) (in a complete agreement with the above arguments).

The purpose is to construct the effective action in the low-energy limit, i.e. below the thresholds for all massive particles. The contributions of massive particles into the running coupling constant are fixed at the masses and, therefore, their contributions to the parameter \(z\) will be proportional to \(M/m\) in a power, defined by the corresponding coefficient of the \(\beta\)-function. It is much more difficult to investigate the contributions of massless gauge fields. Of course, in this case the coupling constant is not fixed at a definite value and we need to perform a detailed analysis of the IR behavior of the theory. Note, that we are interested not in the renormgroup functions but in the effective action, that can be calculated, for example, in the constant field limit. Therefore, the contribution of the massless gauge field to the parameter \(z\) will be a function of these fields, instead of masses.

Because \(z\) is a scalar, in the constant field limit it is necessary to find the chiral scalar superfield, which does not contain the derivatives of \(F_{\mu\nu}\) and anticommuting fields in the lowest component (otherwise all sufficiently large powers of \(z\) will be equal to 0 or infinity). The only superfield \(B\) satisfying these requirements is defined as

\[
B = -\frac{1}{8} \bar{D}(1 - \gamma_5)D(W_{a}^*)^2 = (D_a)^2 - \frac{1}{2} (F_{\mu\nu}^a)^2 - \frac{i}{2} F_{\mu\nu}^a \tilde{F}_{\mu\nu}^a + O(\theta),
\]

where the index \(a\) runs over the generators of a gauge group. (At the perturbative level similar expression was proposed in \([14]\).)

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where the index \(a\) runs over the generators of a gauge group. (At the perturbative level similar expression was proposed in \([14]\).)
Taking into account dimensional arguments, we find, that the contributions of massless gauge fields to the parameter $z$ should be proportional to $M/B^{1/4}$ in a degree, defined by the corresponding coefficient of the $\beta$-function.

The one-loop $\beta$-function can be presented as

$$
\beta(e) = -\frac{e^3}{16\pi^2}(c_{\text{gauge}} + c_{\lambda} + c_q + c_{sq}),
$$

where we denote

$$
c_{\text{gauge}} = \frac{11}{3}N_c, \quad c_{\lambda} = -\frac{2}{3}N_c, \quad c_q = -\frac{2}{3}N_f, \quad c_{sq} = -\frac{1}{3}N_f
$$

the contribution of Yang-Mills fields (with ghosts), their spinor superpartners, quarks and squarks respectively. Then, taking into account the above arguments, we obtain

$$
z = e^{-8\pi^2/e^2}M^{3N_c-N_f}\left(\frac{m_\lambda}{B^{11/12}}\right)^{N_c}\left(\frac{m_q^{2/3}m_{sq}^{1/3}}{B^{1/12}}\right)^{N_f},
$$

where $m_\lambda$ is a gluino mass, $(m_q)^i_j$ and $(m_{sq})^i_j$ are mass matrixes of quarks and squarks.

It should be noted, that the equation (16) was found in the one-loop approximation. Multiloop effects can be included into the Wilson effective action [13, 10, 11], if we take into account, that the instanton measure [13] contains the factor $(1/e^2)^N_c$, which, certainly, will be present in $z$. Therefore, the parameter $z$ (up to a constant $C$) is finally written as

$$
z = C\left(\frac{1}{e^2}\right)^{N_c}e^{-8\pi^2/e^2}M^{3N_c-N_f}\left(\frac{m_\lambda}{B^{11/12}}\right)^{N_c}\left(\det(m_q)^i_j\right)^{2/3}\left(\det(m_{sq})^i_j\right)^{1/3},
$$

where we took into account $SU(N_f)$ symmetry of rotation in the flavor space and restored indexes of generations.

Let us denote

$$
\Lambda_c = \left(\frac{1}{e^2}\right)^{3/11}m_\lambda^{2/11}\left(Ce^{-8\pi^2/e^2}M^{3N_c-N_f}\left(\det(m_q)^i_j\right)^{2/3}\left(\det(m_{sq})^i_j\right)^{1/3}\right)^{3/(11N_c)}
$$

Then the expression for $z$ can be written in the following form:

$$
z = \left(\frac{\Lambda_c}{B^{1/4}}\right)^{11N_c/3}.
$$

As we will see below, $\Lambda_c$ is a typical scale of confinement.
3 Confinement mechanism

Now, let us note, that the field $D$, although being auxiliary, is a real quantum field and should be integrated out. In the considered low energy limit it is necessary to substitute all squarks wave functions by their vacuum expectation values, which are equal to 0 because the gauge group $SU(3)$ is not broken. Therefore, the equation of motion for the field $D$ should be obtained from the action

$$S_W = \frac{1}{16\pi}\text{Im tr} \int d^4x \, d^2\theta \, W^2 \tau(z^{-1/4}). \quad (20)$$

Then, let us note that due to the supersymmetry and \cite{13} the field $D$ is present in $S_W$ only in the following combination

$$b = (D^a)^2 - \frac{1}{2}(F^a_{\mu\nu})^2 - \frac{i}{2}F^a_{\mu\nu}\tilde{F}^a_{\mu\nu}. \quad (21)$$

Therefore, as can be easily seen performing the integration over the anticommuting variable, the bosonic part of the Lagrangian can be presented as

$$L_{\text{bose}} = \frac{1}{8\pi}\text{Im} \left(b^* f(b)\right), \quad (22)$$

where $f(b) \equiv \tau(z^{-1/4}(b))$.

In the vacuum state the field $D$ has such value, that $L_{\text{bose}}$ is minimal. If the magnetic field is absent, then $\tilde{F}_{\mu\nu}F_{\mu\nu} = 0$ and the parameter $b$ can be considered to be real.

The plot of the function $\text{Im} \left(bf(b)\right)$ is presented at Fig.\ref{fig1}. (For simplicity we set $\Lambda_c = 1$ at all figures.)

Using expressions for $a(u)$ and $a_D(u)$ in terms of elliptic functions \cite{16}, it is easy to see, that $\text{Im} \left(bf(b)\right)$ has the only extremum $b_0$, satisfying the condition

$$a(b_0) = z^{-1/4}(b_0) = 4/\pi \quad (23)$$

In the minimum point $\tau(z^{-1/4}(b_0)) = 0$, so that there is no kinetic term for the gauge field in the vacuum state. Therefore, $A_\mu^a$ is a Lagrange multiplier, which corresponds to the constraint

$$J_\mu^a = 0. \quad (24)$$

Condition \cite{24} means, that there are no asymptotic color state and, therefore, leads to confinement of color charges.

At the first sight from the above consideration one can mistakenly conclude, that color charges can not exist in different space points. Really, the kinetic term for the Yang-Mills theory is usually written as
\[- \frac{1}{4e^2} \text{tr} F_{\mu\nu}^2,\]

where \( F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu + i[A_\mu, A_\nu] \). If the corresponding term is absent, the effective charge becomes infinite, that, in turn, leads to the infinitely strong interaction.

However, investigating the quarks interaction it is necessary to take into account the cromomagnetic field, produced by quarks, in the simplest case

\[ A_\mu^a = \frac{q_a^\mu}{r} \delta_{\mu0} \]

Because in the absence of magnetic field (21) is real and can be written as

\[ b = (D^a)^2 + (E^a)^2, \]

the second term is larger than \( b_0 \) for sufficiently small \( r \). Therefore, it is impossible to reach the minimal value of the function \( \text{Im}(bf(b)) \) by any choice of the real field \( D^a \). As it is seen from the presented plot, the function \( \text{Im}(bf(b)) \) is monotonically growing, so that its minimum corresponds to the minimal value of \( b \), which, certainly, corresponds to \( D = 0 \). Because in this case evidently \( b \neq b_0 \), \( \tau(b) \neq 0 \) and, therefore, gauge field has a kinetic term and the theory becomes the usual quantum cromodynamics.

Thus, there are two quite different phases in the considered model. In the first one, corresponding to \( b_0 > (E_a)^2 \), the value of the auxiliary field \( D \) is not equal to 0 (to be exact \( (D_a)^2 + (E_a)^2 = b_0 \)), there is no kinetic term for the gauge field and color charges are confined. In the second phase \( b_0 < (E_a)^2 \), \( D = 0 \) and the theory can be described by usual methods. The point of ”phase transition” \( r_c \) corresponds to

\[ b_0 = (E_a)^2 \approx \frac{1}{r^4} \]

From (19) and (23) we conclude, that

\[ (b_0)^{1/4} = \Lambda_c \left( \frac{4}{\pi} \right)^{12/(11N_c)}, \]

so that \( \Lambda_c \) is really a typical scale of confinement.

Nevertheless, the above picture does not predict the linear growing of the interquark potential. Instead of it we obtain, that for \( r \), larger than a critical value \( r_c \), the potential is infinite. At the small \( r \), using the motion equations, we obtain that

\[ \nabla \left( \text{Im} f_1(1/r^4) \nabla U \right) = \text{const} \delta(r) \]

where \( f_1(b) = f(b) + bf'(b) \). Therefore, the potential is given by

\[ U(r) \sim \int \frac{dr}{r^2 \left( \text{Im} f_1(1/r^4) \right)} \]
that is actually a Colomn potential, modified by the quantum corrections.

The plot of potential (31) is presented at Fig. 2. (As earlier, $\Lambda_c = 1$ and, moreover, we omit quark charges for simplicity, that gives $r_c = (b_0)^{-1/4}$.) A normalization constant is chosen so that $U(r_c) = 0$.

4 Confinement scale

In this section we will investigate only MSSM QCD, which is a particular case of the considered model corresponding to $N_c = 3$ and $N_f = 6$.

In principle, due to the renorminvariance, $\Lambda_c$ can be calculated at any scale, for example, at the Grand Unification scale. In this case the value of the coupling constant in MSSM should be set to be $e^2 \approx 1/2$ and for the confinement scale from (18) we obtain

$$\Lambda_c = M e^{-16\pi^2/11} \left( \frac{m_q^4 m_s^2 \lambda m_{sq}}{M^8} \right)^{1/11}.$$  (32)

Unfortunately, we do not know masses of superpartners. Moreover, in this expression we omit a constant factor, which can be rather large and slightly change the result. (From the other side, the corresponding contribution is close to 1, because this factor is in 1/11 degree).

Nevertheless, a rough estimate of $\Lambda_c$ can be made. Setting all masses equal to 10 – 100 GeV ($M = 2 \times 10^{16}$ GeV), we obtain that

$$\Lambda_c \approx 0.09 - 0.46 \text{ GeV}$$  (33)

that is in a good agreement with the experimental data.

5 Conclusion

In the present paper we propose a mechanism of confinement, which is quite different from the dual Higgs mechanism [5], that is usually used for understanding of this phenomenon. However, this mechanism is a rather natural consequence of the structure of exact effective action (11), with the parameter $z$, given by equation (17).

However, it should be noted, that (14) is only a hypotheses, that can be confirmed (rejected or modified) only by the perturbative and instanton calculations, similar to [17]. In the paper we presented some facts in its favor, for example, the agreement with the exact Novikov, Shifman, Vainshtein and Zahkarov $\beta$-function [10] (that can be easily verified) or the positiveness of the effective charge. Nevertheless, they can not be considered as a strict proof.
In the mechanism of confinement, proposed in this paper, the two main points are the most essential: First, it is the existence of auxiliary field $D$ in supersymmetric gauge theories. Second, we use highly nontrivial form of the bosonic part of the effective action, which is obtained by summing a series of instanton corrections.

However, the interquark potential is appeared not to have linear growing at the large distances. Instead of it there are two different regions (or "phases"). In the one phase (at distances, larger than a critical one) the kinetic term for the gauge field is absent and the potential is infinitely large. In the other phase the theory behaves in a standard way, the potential being approximately equal to the Colomn one. (Similar potentials are used in the bag models [18]).

Nevertheless, the proposed mechanism of confinement is not reduced to the existence of this potential. Actually we predict the absence of color states at the large distances, that is actually observed in the nature. Moreover, we automatically obtain the confinement of electric (instead of magnetic) charges and need not to use duality.

The confinement scale can be calculated from the first principles and is in a good agreement with the experiment, especially taking into account the absence of experimental data on the superpartners masses.

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References

[1] Simonov Yu.//Phys.Usp. 1996. V.39. P.313. (e-print hep-ph/9709344)
[2] Seiberg N., Witten E.//Nucl.Phys. 1994. V.B426. P.19.
[3] Review of Particle Properties//Phys.Rev. 1994. V.D50. P.1173.
[4] Dienes K. e-print hep-th/9602043.
[5] Mandelstam S.//Phys.Lett. 1975. V.53B. P.476;
’t Hooft G.//Nucl.Phys. 1981. V.190. P.455.
[6] ’t Hooft G.//Phys.Rev.Lett. 1976. V.37. P.8;
’t Hooft G.//Phys.Rev. 1976. V.D14. P.3432.
[7] Affleck I., Dine M., Seiberg N.//Nucl.Phys. 1984. V.B241. P.493.
[8] Seiberg N.//Phys.Rev. 1994. V.D49. P.6857.
[9] Pronin P., Stepanyantz K.//Teor.Mat.Fiz. 1999. V. 120. P.82. (e-print hep-th/9902163).

[10] Novikov V., Shifman M., Vainstein A., Zakharov V.//Nucl.Phys. 1983. V.B229. P.381.

[11] Shifman M., Vainshtein A. e-print hep-th/9902018.

[12] Bonelli G., Matone M., Tonin M.//Phys. Rev. 1997. V.D55. P.6466; 
    Flume R., Magro M., O’Raifeartaigh L., Sachs I., Schnetz O.//Nucl.Phys. 1997. 
    V.B494. P.331.

[13] Cordes S.//Nucl.Phys. 1986. V.B273. P.629.

[14] Novikov V., Shifman M., Vainshtein A., Zakharov V.//Phys.Lett. 1986. V.166B. 
    P.329.

[15] Shifman M., Vainshtein.//Nucl.Phys. 1986. V.B277. P.456.

[16] Alvares-Gaume L., Hassan S.//Fortsch.Phys. 1997. V.45. P.159.

[17] Yung A.//Nucl.Phys. 1997. V.B485. P.38.

[18] Close F. An introduction to quarks and partons. Academic Press. London New 
    York San Francisco. 1979.
Figure 1: The plot of function $\text{Im}(bf(b))$ (curve 1). For the comparisons we present the plot (curve 2) of the corresponding perturbative expression.
Figure 2: The plot of interquark potential $U(r)$ in MSSM QCD.