Siberian Branch of Russian Academy of Science
BUDKER INSTITUTE OF NUCLEAR PHYSICS

V.L. Chernyak

ON PROPERTIES OF $\mathcal{N} = 1$ SUSY YANG-MILLS VACUUMS AND DOMAIN WALLS

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The properties of $\mathcal{N}=1$
Susy Yang-Mills vacuums and domain walls

Victor Chernyak

Budker Institute of Nuclear Physics
630090 Novosibirsk, Russia

Abstract

It is shown that there is no chirally symmetric vacuum state in the $\mathcal{N}=1$ supersymmetric Yang-Mills theory. The values of the gluino condensate and the vacuum energy density are found out through a direct instanton calculation. A qualitative picture of domain wall properties is presented, and a new explanation of the phenomenon of strings ending on the wall is proposed.

e-mail: chernyak@inp.nsk.su

1. The $\mathcal{N}=1$ supersymmetric Yang-Mills theory (SYM) partition function is (Q is the topological charge):

$$Z=\sum_k \int dA_\mu d\lambda d\bar{\lambda}\,\delta(Q-k)\,\exp\left\{\frac{i}{4g_o^2}\int dx d\theta\,\frac{1}{2}\,W_\alpha^2+h.c.\right\}, \tag{1}$$

Let us integrate it now over the gluon and gluino fields, but with the chiral superfield W^2-fixed. Proceeding as in [1], one obtains the partition function for the chiral superfield $\Omega=(W^2/32\pi^2N_c)$ in the form:

$$Z=\sum_k \int d\Omega\, d\bar{\Omega}\,\delta(Q-k)\,\exp\{i\int dx L\},$$

$$\frac{1}{N_c^2}L=\left\{\frac{1}{2}\int d\theta\,\Omega\ln\left(\frac{\Omega}{e\Lambda_N^3}\right)+h.c.\right\}+\int d\theta d\bar{\theta}\,M(\Omega,\overline{\Omega},D^n\Omega,\overline{D^n\Omega},\ldots). \tag{2}$$

The F-term in eq.(2) which accounts for all (super) anomalies coinsides with the well known Veneziano-Yankielowitz (VY) [2] effective Lagrangean. The difference is that its real origin and meaning (see [1] for more detail) are clear here from its derivation, while the meaning of the word "effective" was not quite clear for the VY Lagrangean, as well its connection with the original fundamental Lagrangean.

The D-term in eq.(2) is nonanomalous and depends both on the field Ω and its superderivatives. (For our purposes, we will ignore in what follows all fermionic components of Ω and all terms with usual space-time derivatives).

We will show in this section that the chirally symmetric vacuum state obtained by A.Kovner and M.Shifman [3] (KS-vacuum with $\langle 0|\lambda\lambda|0\rangle=0$) is an artefact of using the total VY-Lagrangean, i.e. with the D-term in eq.(2) chosen in the form:

$$M=const\;(\bar{\Omega}\Omega)^{1/3}. \tag{3}$$

In what follows, we prefer to deal with the usual component fields:

$$\Omega=(\sigma,\,\theta^2\chi),\quad \sigma=\frac{\lambda\lambda}{32\pi^2N_c},\quad \chi=S+iP=\frac{GG+iG\tilde{G}}{32\pi^2N_c}, \tag{4}$$

so that the VY-potential takes the form:

$$\frac{1}{N_c^2}U = \frac{1}{2}\left\{(S+iP)\ln\left(\frac{\sigma}{\Lambda^3}\right) + h.c.\right\} - C_o\frac{S^2+P^2}{|\sigma|^{4/3}} \,. \tag{5}$$

With this form, there are N_c chirally asymmetric vacuum states:

$$\bar{\sigma}_n = \langle 0|\sigma|0\rangle_n \sim \langle 0|\lambda\lambda|0\rangle_n \sim \Lambda^3 \exp\left\{i\frac{2\pi n}{N_c}\right\},\ n = 0, ..., N_c - 1, \tag{6}$$

corresponding to the spontaneously broken residual axial symmetry and besides, as emphasized by A.Kovner and M.Shifman [3], there is also the chirally symmetric vacuum solution:

$$\langle 0|\lambda\lambda|0\rangle_o = 0. \tag{7}$$

Let us point out first that two solutions, eq.(6) and eq.(7), are not on equal footing. Because we know (from the Witten index) that SUSY is unbroken, we are ensured that $\overline{S} = \langle 0|S|0\rangle \to 0$. So, if $|\overline{\sigma}| \neq 0$, it is sufficient to use eq.(5) to find out the value of $\overline{\sigma}$, as higher order terms, like $S(S/|\sigma|^{4/3})^k$, are of no importance in this case. If $|\overline{\sigma}| \to 0$ however, all higher order terms become of importance and we can not believe, in general, the results obtained from eq.(5). If the VY-potential were exact, the KS-solution will survive. But really, the term $S^2/|\sigma|^{4/3}$ in eq.(5) is only the first term in the expansion in powers of $(S/|\sigma|^{4/3})^k$. So, the KS-solution is not selfconsisted in this respect and we need to know, in particular, the behaviour of the potential at $z = S/|\sigma|^{4/3} \to \infty$.

To find it out, let us write first a general form of the potential in eq.(2) ($\sigma = \rho\exp\{i\phi\}$):

$$\frac{1}{N_c^2}U(\sigma,\chi) = S\ln\left(\frac{\rho}{\Lambda^3}\right) - (\phi - \frac{\theta}{N_c})P + Sf\left(\frac{S}{\rho^{4/3}}, \frac{P}{S}\right). \tag{8}$$

Let us add now to eq.(8) the gluino mass term:

$$\frac{1}{N_c^2}\Delta U = -m_o\,\rho\cos(\phi)\,, \tag{9}$$

where m_o is the renormalization group invariant mass parameter. [1] This addition of ΔU is legitimate as our Lagrangian was obtained integrating out all degrees of freedom, but with all components of the Ω-superfield fixed.

[1] Changing the phase of m_o is equivalent to a redefinition of θ in eq.(8). So, it is convenient to choose m_o in eq.(9) to be real and positive.

Now, at large $m_o \to \infty$, the heavy gluino will decouple leaving us with the pure YM theory and we know how it decouples, from the renormalization group. In this region (see below): $\overline{S} = O(m_o^{8/11})$, $\overline{\rho} = O(m_o^{-3/11})$, so that $\overline{S}/\overline{\rho}^{4/3} = O(m_o^{12/11}) \to \infty$. Therefore, this will allow us to find out the asymptotic behaviour of f in eq.(8).

As the gluino contribution to $b_o = 3 = (11/3 - 2/3)$ is $(-2/3)$, it is not difficult to check that the function f in eq.(8) has to have the asymptotic behaviour:

$$f\left(\frac{S}{\rho^{4/3}}, \frac{P}{S}\right) \to \frac{1}{4}\ln\left(\frac{S}{\rho^{4/3}}\right) + \tilde{f}\left(\frac{P}{S}\right),\quad \frac{S}{\rho^{4/3}} \to \infty. \tag{10}$$

In this case, integrating out the ρ and ϕ fields, one has:

$$m_o\overline{\rho}\,e^{i\overline{\phi}} = \left(\frac{2}{3}S + iP\right), \tag{11}$$

and $U(S,P)$ has the form:

$$\frac{1}{N_c^2}U(S,P) = \frac{11}{12}S\ln\left(\frac{S}{\Lambda_{YM}^4}\right) + \frac{\theta}{N_c}P + S\tilde{f}\left(\frac{P}{S}\right);\ \Lambda_{YM} = \Lambda^{9/11}m_o^{2/11}, \tag{12}$$

as it should be. [2]

Now, we are ready to check the existence of the KS-solution: $\overline{\rho} \to 0$. We can distinguish three cases (we take $\theta = 0$, $\overline{\phi} = \overline{P} = 0$, as they are of no importance for us here).

a) Let $\overline{z} = (\overline{S}/\overline{\rho}^{4/3}) \to 0$, so that $f(z = S/\rho^{4/3}) \sim z$. As was pointed out above, this variant is selfcontradictory at $\overline{\rho} \to 0$, as $\partial U/\partial S = 0$ leads to: $\overline{z} \sim \ln(\Lambda^3/\overline{\rho}) \to \infty$.

b) Let $\overline{z} \to z_o = const \neq 0$. Then (barring pathological singularities) the saddle point equations are: $z_o f'(z_o) = 3/4$; $\ln(\Lambda^3/\overline{\rho}) = f(z_o) + z_o f'(z_o)$. The first equation shows that $f'(z)$ (and so $f(z)$) is nonsingular at $z = z_o$, but we are in trouble then with the second equation at $\overline{\rho} \to 0$.

c) Finally, let $\overline{z} \to \infty$, so that $f(z) \to (1/4)\ln z$. This case is also in trouble, as $\partial U/\partial S = 0$ leads to $(\overline{S}/\overline{\rho}) = O(1/\overline{\rho}^{11/3}) \to \infty$ at $\overline{\rho} \to 0$, while $\partial U/\partial\rho = 0$ leads to $(\overline{S}/\overline{\rho}) \to 0$.

On the whole, we conclude that there is no chirally symmetric vacuum state in $\mathcal{N} = 1$ SYM, so that the residual axial symmetry is spontaneously broken in all vacuum states.

[2] Another way to check eq.(10) is to recall that eq.(11) can be obtained through a direct calculation of the heavy gluino loop in the gluon background, and is directly related to the trace and axial anomalies.

2. We will show now that a spontaneous breaking of the residual axial symmetry and the value of the gluino condensate can be obtained in a quite different way, through a direct calculation of the instanton contributions into the partition function. With this purpose, let us return to the original partition function, eq.(1), add the gluino mass term with a small but finite mass m_o to the action, and consider the instanton contributions.

It has been shown in [4] that, under a special choice of the collective coordinates, the n-instanton contribution splits up into nN_c "instantonic quarks". In our case of $\mathcal{N}=1$ SYM, the result is especially simple. Because all nonzero mode contributions cancel exactly between the gluon and gluino contributions, there remains no residual interaction between these instantonic quarks. For instance, the n=1 instanton contribution takes the form ($b_o=3$, $m_\theta = m_o \exp{(i\,\theta/N_c)}$):

$$Z_1 = \int dx_1 ... dx_{N_c} \frac{1}{N_c!}\left[N_c^2 \frac{m_\theta}{2} \Lambda^{b_o}\right]^{N_c}, \qquad (13)$$

where: x_i are the collective coordinates (the centre positions of the instantonic quarks), and m_θ is due to the gluino zero modes. The n-instanton contribution is exactly of the same form, so that summing up over n (and adding antiinstantons) one obtains the partition function in the form:

$$Z_{tot} = Z\,Z^*, \quad Z = \frac{1}{N_c}\sum_{k=0}^{N_c-1} e^{I(k)},$$

$$I(k) = \int dx\, N_c^2 \left\{\frac{m_\theta}{2}\Lambda^3\left[1+O(|m_\theta|^2)\right]\exp\left(i\frac{2\pi k}{N_c}\right) + O(|m_\theta|^2)\right\}. \quad (14)$$

Here, the factor $Z_{N_c}(k) = \exp\{i\,2\pi k/N_c\}$ appeared because we have extracted the N_c-th power root from unity, when going from eq.(13) to eq.(14). It plays the role of the "neutralizator", i.e. when $\exp\{I(k)\}$ in eq.(14) is expanded back into a power series, it ensures that instantonic quarks appear in the N_c-fold clusters only (i.e. in the form of instantons). The same, it ensures the periodicity: $Z(\theta) = Z(\theta + 2\pi l)$, which was explicit before summation over n.

We would like to emphasize that the above expression for the action in eq. (14) is exact, within the indicated accuracy. Indeed:

a) The perturbation theory (i.e. the $Q=0$ sector of the partition function) contribution is exactly zero at $m_o = 0$ due to SUSY, and the corrections from this sector start with $O(|m_\theta|^2)$. This is because the replacement $m_o \to -m_o$ is equivalent to changing the θ- angle, and the $Q=0$-sector is θ-independent.

b) Because the one-loop Z_Q-contribution contains already the factor $(m_\theta)^{QN_c}$, all higher loop corrections to it can be calculated with $m_o=0$, and they all cancel due to SUSY. For the same chirality reasons, the (relative) corrections in this sector also start with $O(|m_\theta|^2)$, including those which originate from disturbing the exact cancelation between the gluon and gluino nonzero modes at $m_o = 0$.

c) As for the instanton-antiinstanton interaction contributions, they should not be considered as independent ones, but rather as belonging to perturbation theory (its asymptotic tail), in the sector with fixed Q. So, they are also zero in the sense that they are accounted for already in the points "a" and "b" above.

Eq.(14) shows clearly that the residual axial symmetry is broken spontaneously in our system (in the infinit volume limit). Indeed, before summation over n each n-instanton contribution was invariant by itself under $\theta \to \theta + 2\pi l$, as a result of the residual axial symmetry. But after summation, the instantonic quarks have released and the above symmetry acts nontrivially now, interchanging N_c branches between themselves. As a result, because the small perturbation ($m_o \neq 0$) was introduced, one definite branch dominates the whole partition function, - those one which minimizes the energy (at given θ). So, we obtain for the vacuum energy density:

$$E_{vac} = -N_c^2\,\Lambda^3\,\frac{1}{2}\left[\,m_\theta + \overline{m}_\theta\,\right]_{2\pi} + O(|m_\theta|^2), \qquad (15)$$

where the notation $[f(\theta)]_{2\pi}$ means that this function is $f(\theta)$ at $-\pi \le \theta \le \pi$, and is glued then to be periodic in $\theta \to \theta + 2k\pi$, i.e.: $[f(\theta)]_{2\pi} = \min_k f(\theta + 2\pi k)$.

Further, because the $O(m_o)$ term appeared in the energy, this shows that the order parameter (the gluino condensate) is nonzero. Indeed, let us consider [3] :

$$V^{-1}\exp\left\{\frac{i\theta}{N_c}\right\}\left[\frac{\partial \ln Z}{\partial(m_\theta/2)}\right]_{m_o=0} = \sum_{k=0}^{N_c-1}\langle\theta, k|\lambda\lambda|\theta, k\rangle =$$

$$= \sum_{k=0}^{N_c-1} N_c\,\Lambda^3 \exp\left\{i\,\frac{\theta + 2\pi k}{N_c}\right\}. \qquad (16)$$

Thus, it is clear from eqs. (14), (16) that (at $m_o \to 0$), there are N_c degenerate vacuum states differing by the phase of the gluino condensate: $\langle\theta, k|\lambda\lambda|\theta, k\rangle = N_c\,\Lambda^3 \exp\{i\,(\theta + 2\pi k)/N_c\}$.

[3] We can keep m_o infinitesimal but finite, and $V \to \infty$, and to separate out one term from the sum over "k" in eq. (16).

Really, it is possible to replace m_o by the local function $m_o(x)$ in eqs. (13) and (14). Indeed, in eq. (13) the gluino zero mode contributions will take the form:

$$I_o = \int dx_1 \dots dx_{N_c} \int dy_1\, m_\theta(y_1) \dots dy_{N_c}\, m_\theta(y_{N_c})\, \Pi\,,$$

$$\Pi = |\psi_o^{(1)}(y_1 - x_k)|^2 \dots |\psi_o^{(N_c)}(y_{N_c} - x_k)|^2,$$

where $\psi)_o^{(i)}(y_i - x_k)$ means $\psi_o^{(i)}(y_i - x_1, y_i - x_2, \dots, y_i - x_{N_c})$, and $\int dy|\psi_o^{(i)}(y - x_k)|^2 = 1$. It is not difficult to see that: $I_1 = \int dx_1 \dots dx_{N_c}\, \Pi = 1$. Indeed, let us take temporarily $m_\theta = 1$ and put our fields into a large 4-dimensional Euclidean volume V, of an arbitrary form. Because $\int dy_1 \dots dy_{N_c}\, \Pi = 1$, $I_o(m = 1) = V^{N_c}$. Now, if I_1 were a nontrivial function of the ratios like $(y_1 - y_2)^2/(y_2 - y_3)^2$ etc., then $I_o(m = 1)$ will be of the form: $I_o(m = 1) = V^{N_c} f_{geom}$, where the function f_{geom} will depend on the geometry of our volume, and this will be a wrong answer. Therefore, $I_o = \prod_{i=1}^{N_c} \int dy_i\, m_\theta(y_i)$, and m_θ can be replaced by $m_\theta(x)$ in eqs. (14) and (15). While the corrections $O(|m_\theta|^2)$ in eq. (14) remain uncontrolable, it is important that there are no uncontrolable pure chiral corrections of the type $O(m_\theta^l)$, $l \geq 2$. As a result, taking derivatives $\sim \delta/\delta\, m_\theta(x)$ we can obtain even local pure chiral Green functions, like: $\langle k|\lambda\lambda(x_1) \dots \lambda\lambda(x_l)|k\rangle$, and all of them will be pure constants. [4]

There is nothing mysterious in this behaviour and it does not imply that theory is trivial. For instance, let us consider $\langle k|\lambda\lambda(x)\, \lambda\lambda(0)|k\rangle$, and let us denote: $\lambda\lambda(x) = \exp(i\, 2\pi k/N_c)\, \rho(x) \exp(i\phi_k(x))$, so that: $\langle k|\rho(x)|k\rangle = N_c \Lambda^3$ and $\langle k|\phi_k(x)|k\rangle = 0$. Then ($\rho \exp(i\phi_k) = \sigma_k + i\pi_k$):

$$\langle k|\lambda\lambda(x)\, \lambda\lambda(0)|k\rangle = \langle k|\lambda\lambda(x)|k\rangle \langle k|\lambda\lambda(0)|k\rangle +$$

$$+ \exp\left\{\frac{i\, 4\pi k}{N_c}\right\} \{\langle k|\sigma_k(x)\sigma_k(0)|k\rangle_{con} - \langle k|\pi_k(x)\pi_k(0)|k\rangle_{con}\} = \langle k|\lambda\lambda(0)|k\rangle^2,$$

as the nontrivial connected correlators cancel each other due to SUSY.

As was pointed out above, supposing only that the gluino condensate is really nonzero, it becomes legitimate to use the VY Lagrangian to investigate the vacuum properties, i.e. to find out the gluino condensate, eq.(6) [2], and the vacuum energy density, eq.(15) [5]. In other words, it is not an approximation in this case as higher order terms are of no importance. In

[4] Another way to obtain the same result is to introduce a small local mass into eq.(9): $m_o \to m_o(x)$.

contrast, if we want to deal with some exitations, say domain walls, the VY Lagrangean is insufficient.

3. Because there is a spontaneous breaking of the residual axial symmetry, there are the domail wall excitations interpolating between the above N_c chirally asymmetric vacuums. The purpose of this section is to give a new qualitative description and interpretation of the domain wall properties and, in particular, their ability to screen the quark charge.

Let us recall in short the interpretation of the vacuum energy density behaviour in the pure gluodynamics, which was proposed in [1] (we take $N_c = 2$ for simplicity as the qualitative picture at $N_c > 2$ remains the same in essence, and only the number of independent dyon species increases, see [1]).

The vacuum state at $\theta = 0$ is supposed to be the condensate of pure magnetic monopoles, i.e. of the dyons with the magnetic and electric charges $d_1^{\theta=0} = (1, 0)$. As has been shown by E.Witten long ago [6], as θ becomes nonzero the above monopoles turn into the dyons with the charges: $d_1^\theta = (1, \theta/2\pi)$. For this reason, the vacuum energy density, $\bar{U}(\theta)$, increases. This continues up to $\theta \to \pi$ where the above dyons look as: $d_1^{\theta=\pi} = (1, 1/2)$. The above vacuum becomes unstable in the infinitesimal vicinity of $\theta = \pi$ because there is another state, the condensate of $d_2^{\theta=\pi} = (1, -1/2)$ - dyons, degenerate in energy with the first one. Thus, there occurs rearragement of the electrically charged degrees of freedom to recharge the d_1 - dyons into the d_2 - ones. For instance, a copious production of the charged gluons, $\bar{g} = (0, -1)$ takes place, so that: $(\bar{g}) + d_1 \to d_2$. This rechargement allows the system to have a lower energy at $\theta > \pi$. Indeed, there are now only the $d_2^\theta = (1, -1 + \theta/2\pi)$ - dyons in the condensate at $\theta > \pi$, their electric charge decreases with increasing θ and the vacuum energy density decreases with it. As $\theta \to 2\pi$, the d_2^θ - dyons become the pure monopoles, and the vacuum state becomes exactly as it was at $\theta = 0$. On the whole, the vacuum energy density, $\bar{U}(\theta)$, increases in some way at $0 \leq \theta \leq \pi$; there is a cusp due to the above described rechargement at $\theta = \pi$, and it decreases then (in a symmetric way) reaching its minimal value at $\theta = 2\pi$.

Now, let us return to SYM and let us suppose that we have integrated out all, but the ρ and ϕ ($\lambda\lambda \sim \rho \exp\{i\phi\}$) fields (really, we expect the field ρ is unimportant for a qualitative picture discussed below and we will ignore it, supposing simply that it takes its vacuum value $\bar{\rho} = \Lambda^3$).

As the field $N_c\phi$ in SYM is the exact analog of θ in YM, the above described interpretation of the behaviour of $\bar{U}(\theta)$ in YM can be transfered to SYM, with only some evident changes: a) $U(\phi)$ is not the vacuum energy density now but rather the potential of the field ϕ; b) if we start with the condensate of pure monopoles at $\phi = 0$, the rechargement $d_1^\phi = (1, 2\phi/2\pi) \rightarrow d_2^\phi = (1, -1 + 2\phi/2\pi)$ and the cusp in $U(\phi)$ will occur now at $\phi = \pi/2$ for $N_c = 2$ (and at $\phi = \pi/N_c$ at $N_c > 2$), so that at $\phi = \pi$ we will arrive at the next vacuum with the same pure monopole condensate but with the shifted phase of the gluino condensate.

Let us consider now the domain wall exitation, $\phi_{dw}(z)$, interpolating along the "z" axis between, say, two nearest vacuums: $\phi(z \rightarrow -\infty) \rightarrow 0$ and $\phi(z \rightarrow \infty) \rightarrow \pi$. There is a crucial difference between this case and those just described above where the field ϕ was considered as being space-time independent, i.e. $\phi(z) = const$. The matter is that the system can not behave now in a way described above (which allowed it to have a lowest energy at each given value of $\phi(z) = \phi = const$): i.e. to be the pure condensate of d_1^ϕ - dyons at $0 \leq \phi < \pi/2$, the pure condensate of d_2^ϕ - dyons at $\pi/2 < \phi \leq \pi$, and to recharge suddenly at $\phi = \pi/2$). The reason is that the fields corresponding to electrically charged degrees of freedom also become functions of "z" at $q = \int dz\,[d\phi_{dw}(z)/dz] \neq 0$. So, they can not change abruptly now at some $z = z_o$ where $\phi_{dw}(z)$ goes through $\pi/2$, because their kinetic energy will become infinitely large in this case. So, the transition will be smeared necessarily.

Therefore, there will be some admixture of the d_2^ϕ - dyons even in the $0 < \phi(z) \leq \pi/2$ region, and of the d_1^ϕ - dyons in the $\pi/2 \leq \phi < \pi$ region. If we denote the fractions of our $d_{1,2}$-dyons inside the domain wall by $x_{1,2}(z)$, $x_1 + x_2 = 1$, [5] then, for instanse: $x_1(z) \rightarrow 1$ at $z \rightarrow -\infty$, $x_1(z_o) = x_2(z_o) = 1/2$, and $x_1(z) \rightarrow 0$ at $z \rightarrow \infty$.

Finally, let us consider what happens when a heavy quark is put inside the domain wall. The crucial point is that there is the nonzero mixture of four dyon and antidyon species: $d_1^\phi = (1, 2\phi/2\pi)$, $\bar{d}_1^\phi = (-1, -2\phi/2\pi)$, $d_2^\phi = (1, -1 + 2\phi/2\pi)$ and $\bar{d}_2^\phi = (-1, 1 - 2\phi/2\pi)$, through all the domain wall profile. Polarizing itself appropriately, this mixture of dyons is capable to screen any test charge which has been put inside (see [1] for more detail). And as the string is not developed, the energy is lowered in this way. If the test quark is put at far left (right), $z \rightarrow \pm\infty$, the string will originate from this point making its way toward a wall. But on the way, it will be more and more screened and will disappear finally in the bulk of the wall. This explanation differs from both, those described by E.Witten in [7] and those proposed by I.Kogan, A.Kovner and M.Shifman in [8].

[5] There may be some admixture of higher excited states, but we will ignore them as they do not change the qualitative picture.

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References

[1] V. Chernyak, preprint BINP 98-61, hep-th/9808092

[2] G. Veneziano and S. Yankielowicz, Phys. Lett. **B113** (1982) 321

[3] A. Kovner and M. Shifman, Phys. Rev. **D56** (1997) 2396

[4] V.A.Fateev, I.V.Frolov and A.S.Shvarts, Sov.J.Nucl.Phys. **30**(1979)590
B. Berg and M. Luscher, Comm.Math.Phys. **69** (1979) 57
R.D. Carlitz and D.A. Nicole, Nucl. Phys. **B243** (1984) 307

[5] A. Masiero and G. Veneziano, Nucl. Phys. **B249** (1985) 593

[6] E. Witten, Phys. Lett. **B86** (1979) 283

[7] E. Witten, Nucl. Phys. **B507** (1997) 658; hep-th/9706109

[8] I.I. Kogan, A. Kovner and M. Shifman, hep-th/9712046

Victor Chernyak

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