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THE DISCRETE CHIRAL SYMMETRY BREAKING IN QCD AS A MANIFESTATION OF RYBAKOV-CALLAN EFFECT

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THE DISCRETE CHIRAL SYMMETRY BREAKING IN QCD AS A MANIFESTATION OF RYBAKOV-CALLAN EFFECT

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Abstract

In the presence of topological monopole-type fluctuations (merons), the discrete chiral symmetry breaking takes place. This phenomenon resembles the Rubakov-Callan effect. The connections of this phenomenon with Gribov's ambiguity and with fermion number fractionization are discussed. It is shown that the Gribov's vacuum has the nonzero axial number.

1. In this note we discuss the properties of well-separated meron-antimeron fluctuation [1]. It is well-known [1] that the meron configuration has monopole-type component. So, the effect of Rubakov-Callan-type [2,3] (the chiral condensate appears in the presence of magnetic monopoles) can arise. Really, below it will be shown that in the presence of meron's configuration the whole variation of chiral charge (during the life-time of meron) is nonzero.

The another interesting property of meron is the following one. The meron connects [4] trivial and Gribov's vacua [5]. Our main purpose of this note - is the analysis of this Gribov vacuum.

The example of vacuum fluctuation with the finite action, which creates the quasi-Gribov state (weakly dependent on time-variable) is shown at fig. 1. In the following we will neglect this nonstationary.

It will be shown that static Gribov vacuum ($A_0 = 0$, $A_i = -i U^+ \partial_i U$, $U = \exp\{ i g(r) \frac{\vec{n}\vec{\sigma}}{2} \}$, $g(\infty) = \pm\pi$, $g(0) = 0$) is not equivalent to trivial vacuum field ($A_\mu = 0$) in the presence of fermions*. The consequence (and independent proof) of this is existence of normalizable zero-energy eigenmode in the $A_i = -U^+ i \partial_i U$ field. This means that the soliton A_i^a interpolate between the two vacua. Hence the Gribov states are doubly degenerate and they are marked by chiral charge Q_5 . So we have states with discrete chiral symmetry (CS): $|Q_5 = +1\rangle \leftrightarrow |Q_5 = -1\rangle$ being spontaneously broken: $\psi \to \exp\{ i \frac{\pi}{2} \gamma_5 \} \psi$.

Let us note, that in one space dimension, when discrete symmetry is spontaneously broken, there are always solitons, the states interpolating between two different vacua [6]. Because of spherical symmetry of field A_i^a our problem can be reformulated in the 2-d language. And analysis of discrete system of vacua can be carried out analogously to that of ref. [6].

* It is so because $U \to \vec{n}\vec{\sigma}$ when "r" tends to infinity ($U \neq c, r \to \infty$). Hence U is not acceptable as gauge transformation. In other words, U does not satisfy boundary condition ($U(0) = U(L) = c$) in the box, see p. 5.

As mentioned above the meron-antimeron fluctuation have finite contribution to functional integral because normalizable $\int d^4x$ zero energy eigenmode is absent in this case. Hence, the phenomenon of spontaneous breaking of discrete CS, which is discovered in one fluctuation is left in full theory.

In the p. 2 the problem is reformulated into 2-d field theory with some boundary condition (analogously to [2]). The "bosonization method [7] is very useful for analysis of this 2-d theory.

Another independent approach (point 3) to this problem is connected with the calculation of value φ_5 in the presence of Gribov's field. This value reduces to the fermion number of 2d soliton. The suitable methods of calculation of this fermion numbers are well-known [8]. The physical sense of this approach is partially lost (the role of the monopole-configurations, boundary conditions and so on are not seen). However, the reducing of the problem to the calculation of fermion number of corresponding solitons is very useful. Due to this reducing we can reformulate the problem in topological language (point 4), without restriction by spherically-symmetric fields.

2. Before the reformulation of problem in the 2d language, let us discuss some questions, concerning Euclidean formulation of theory and notations.

For the reformulation of theory in Euclidean language we follow paper [9], where in particular it was shown that there is no necessity to change fields $(\psi, \bar{\psi}, A_\mu)$ and γ_μ - matrix for this reformulation. We may carry out calculations in Minkovsky space. However after calculation we have to replace x_0 by $-ix_4$, where $x_4 \equiv t$ is real value. This prescription is not necessary in QCD (in contrast to SUSY-models [9]), but it is very useful.

Taking into account this remark let us consider spherically-symmetric Witten's Ansatz [10] in the following form:

$$W_0^a = n^a A_0 ,$$
$$W_i^a = (\delta^{ai} - n^a n^i)\frac{\varphi_1}{r} + n^a n^i A_1 + \epsilon^{iak} n^k \frac{1+\varphi_2}{r} . \qquad (1)$$

Here $n^i = x^i/r$; $r = \sqrt{x^i x^i}$; $A_\mu (\mu = 0,1)$, φ_1, φ_2 - some functions of x_0, r. In this notation the standart meron's solution $W_\mu = W_\mu^a \frac{\sigma^a}{2} = \frac{1}{2} i\, U^+ i \partial_\mu U$ takes the form:

$$A_0 = \frac{-ir}{(-x^2)}, \quad A_1 = \frac{-ix_0}{(-x^2)}, \quad \varphi_1 = \frac{-ix_0 r}{(-x^2)}, \quad \varphi_2 = \frac{x_0^2}{(-x^2)} \qquad (2)$$
$$-x^2 = -x_0^2 + r^2 = t^2 + r^2 > 0$$

In the $A_0 = 0$ gauge, the solution (2) reduces to the simple form:

$$A_0 = A_1 = \varphi_1 = 0, \quad \varphi_2 = \frac{-ix_0}{\sqrt{-x^2}} . \qquad (3)$$

The topological charge Q

$$Q = -\frac{1}{32\pi^2} \int d^4x\, G_{\mu\nu}^a \tilde{G}_{\mu\nu}^a \qquad (4)$$

is $1/2$ for the configuration (2) and zero for eq. (3) [1].

So, following the papers [2] and take into account S-wave fermions only, let us write Lagrangian L for the effective 2d theory. This Lagrangian describe the spherically symmetric degrees of freedom of QCD with SU(2) gauge group. The classical background field is defined by eq. (3), the quantum fluctuations of vector fields are described by A_μ, $\mu = 0,1$; the S-wave fermions are described by the following form [2]:

$$\psi_1 = \frac{1}{\sqrt{8\pi r^2}} \begin{pmatrix} \varphi_1 \\ -\varphi_1 \end{pmatrix}, \quad \varphi_1^{\alpha i} = \left(g_1 \epsilon^{\alpha i} + i(\vec{n}\vec{\sigma})^{\alpha\gamma} \epsilon^{\gamma i} h_1\right), \quad \chi_1 = \begin{pmatrix} g_1 \\ h_1 \end{pmatrix} \qquad (5)$$

$$\psi_2 = \frac{1}{\sqrt{8\pi r^2}} \begin{pmatrix} \varphi_2 \\ \varphi_2 \end{pmatrix}, \quad \varphi_2^{\alpha i} = \left(g_2 \epsilon^{\alpha i} + i(\vec{n}\vec{\sigma})^{\alpha\gamma} \epsilon^{\gamma i} h_2\right), \quad \chi_2 = \begin{pmatrix} g_2 \\ -h_2 \end{pmatrix}$$

Here $\alpha(i) = 1,2$ are indexes describing spin and color correspondly, $\epsilon^{12} = 1$. In notation (1), (5) we have:

$$S = S_{cl} + \int dr\, dx_0\, L$$

$$L = \bar{\chi}_1 i\hat{\partial}\chi_1 + \bar{\chi}_2 i\hat{\partial}\chi_2 - \frac{\varphi_2(r,x_0)}{r}\left[\bar{\chi}_1 i\gamma_5\chi_1 + \bar{\chi}_2 i\gamma_5\chi_2\right] - \qquad (6)$$
$$-\frac{\varphi_1}{r}\left[\bar{\chi}_1\chi_1 - \bar{\chi}_2\chi_2\right] + \frac{4\pi r^2}{g^2}\left[\frac{\varphi_2^2}{r^2} A_\mu^2 - \frac{1}{4} F_{\mu\nu}^2\right] + \frac{1}{2} A_\mu \left[\bar{\chi}_1\gamma_\mu\gamma_5\chi_1 - \bar{\chi}_2\gamma_\mu\gamma_5\chi_2\right].$$

Here $\varphi_2^{cl.} \neq 0$ only, $F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu$, γ_μ - matrices are that of 2d space:

$$\gamma_\mu \gamma_\nu \{\} = 2 g_{\mu\nu}\,, \quad g_{\mu\nu} = \begin{pmatrix} 1 & \\ & -1 \end{pmatrix}, \quad \gamma_0 = \sigma_3\,, \quad \gamma_1 = -i\sigma_1\,, \quad \gamma_5 = \sigma_2 \tag{7}$$

Note that $\psi_1(\chi_1)$, $\psi_2(\chi_2)$ interact with A_μ, φ_1 fields with the different signs (this fact was noted in ref. [2]). Besides that the parameter $\tilde{m} = \varphi_2/r$ can be understood as effective mass of χ - fields. It is very importqnt that mass $\tilde{m}$ changes sign when time-variable ($t = -\infty \to t = +\infty$) does. This fact leads to the different boundary condition depending on sign of t (8). All this leads to the nonzero variation of value ΔQ_5 when A_μ^a interpolate between the trivial ($A_\mu = 0$) and Gribov ($A_i = -i\,U^+ \partial_i U$) $\underbrace{\text{vacua}}$. Let us show this.

The boundary condition are connected with the requirement of absence of singularities near the point $r = 0$ [2]. Since $\varphi_2 \to 1\,(t<0)$, $\varphi_2 \to -1\,(t>0)$, we have:

$$\begin{aligned} g_{1,2}(r=0) &= 0\,, \quad \sigma^- \chi_{1,2} = 0\,, \quad t < 0, \\ h_{1,2}(r=0) &= 0 \quad \sigma^+ \chi_{1,2} = 0\,, \quad t > 0 \end{aligned} \tag{8}$$

Let us discuss the bose-representation of our 2d theory. The spinor fields χ_1, χ_2 are described by bose-fields (φ_1, φ_2) [7]. Then, massive term $m\bar{\psi}\psi$ in Lagrangian is given by: $\bar{\psi}\psi \sim \bar{\chi}_1\chi_2 + \bar{\chi}_2\chi_1 \sim \cos\sqrt{\pi}(\varphi_2 - \varphi_1)$. We may show that the bose-representation of $\bar{\psi}\psi$ reverses its sign when time - variable does:

$$\begin{aligned} L_m \sim m\bar{\psi}\psi &\sim m\mu \cos\sqrt{\pi}(\varphi_2 - \varphi_1)\,, \quad t > 0 \\ &\sim -m\mu \cos\sqrt{\pi}(\varphi_2 - \varphi_1)\,, \quad t < 0 \end{aligned} \tag{9}$$

Of course this fact is connected with the boundary condition (8). The other terms in Lagrangian are zero for vacua states. That's why, only the term L_m (9) determines the vacuum value $\varphi_2 - \varphi_1$.

On the other hand, the baryon and axial current 4d-theory in the boson-representation have a form:

$$\bar{\psi}\gamma_\mu\psi = \frac{1}{4\pi r^2}\left(\bar{\chi}_1\gamma_\mu\chi_1 + \bar{\chi}_2\gamma_\mu\chi_2\right) = \frac{1}{4\pi r^2}\frac{1}{\sqrt{\pi}}\epsilon_{\mu\nu}\partial_\nu(\varphi_1 + \varphi_2) \tag{10}$$

$$\bar{\psi}\gamma_\mu\gamma_5\psi = \frac{1}{4\pi r^2}\left(\bar{\chi}_1\gamma_\mu\chi_1 - \bar{\chi}_2\gamma_\mu\chi_2\right) = \frac{1}{4\pi r^2\sqrt{\pi}}\epsilon_{\mu\nu}\partial_\nu(\varphi_1 - \varphi_2)\,, \quad \mu = 0, r$$

The whole variation of axial number Q_5 in the presence of meron is equal to:

$$\Delta Q_5 = \lim_{T\to\infty}\left[Q_5(T) - Q_5(-T)\right] = \int \bar{\psi}\gamma_0\gamma_5\psi\, d^3x \Big|_{-T}^{T} = \tag{11}$$

$$\frac{1}{\sqrt{\pi}}\left[\varphi_2(T) - \varphi_1(T)\right]_{r=0}^{r=\infty} - \frac{1}{\sqrt{\pi}}\left[\varphi_2(-T) - \varphi_1(-T)\right]_{r=0}^{r=\infty} = \pm 1$$

In the eq. (11) we take into account the finiteness of Coulomb--interaction at $r = 0$: $\varphi_2(r=0) - \varphi_1(r=0) = 0$. Besides that, the Lagrangian is minimized when $\varphi_2 - \varphi_1 = 0$ for $t < 0$ and $\varphi_2 - \varphi_1 = \pm\sqrt{\pi}$ for $t > 0$ at $r \to \infty$.

So, the meron with half of unit of topological charge connects trivial vacuum and doubly degenerated ($Q_5 = \pm 1$) Gribov state. Hence $Q_5^{\text{Gribov}} = \pm 1$ and $U(1)$ condensate $\langle\bar{\psi}\psi\rangle \neq 0$ is nonzero

In order to understand the role of such configurations we discuss the 2d - theory [11] with analogous property. In ref. [11] it was shown that $\langle\bar{\psi}\psi\rangle \neq 0$ for the 2d theory:

$$L = -\frac{1}{4}F_{\mu\nu}^2 + \bar{\psi}(i\hat{\partial} + e\hat{A})\psi + g^2/N\left[(\bar{\psi}\psi)^2 + (\bar{\psi}\, i\gamma_5\psi)^2\right], \tag{12}$$

due to the fluctuations with topological number $= 1/N$. The result $\langle\bar{\psi}\psi\rangle \neq 0$ means that the theory (12) is noninvariant under the $SU(N)$ and $U(1)$ chiral transformation. The Lagrangian (12) is explicitly noninvariant under chiral $SU(N)$. The $U(1)$ chiral current has an anomalous divergence, $\partial_\mu a_\mu = N\frac{e}{2\pi}\epsilon_{\mu\nu}F_{\mu\nu}$ and does not conserve. But despite the anomaly, a discrete N-fold $CS: \psi \to \exp\{i\pi\gamma_5 k/N\}\psi$ takes place. This discrete symmetry is spontaneously broken due to the fluctuations with nonintegral topological number. This phenomenon provides the nonzero value of $\langle\bar{\psi}\psi\rangle$ [11].

In the case of QCD the situation is quite analogous. The fluctuations with nonintegral topological charge provide the spontaneously broken discrete $U(1)CS$ and can saturate the anomaly. Hence, this fluctuations can influence numerical value of $\langle \bar{\Psi}\Psi \rangle$ if $SU(N)CS$ is already spontaneously broken.

3. Let us discuss more general approach to this problem. By eq. (10) the calculation of fermion number F and axial number Q_5 of 4-d theory reduces to calculation of fermion numbers $F_{1,2}$ of 2-d theory (6):

$$F = F_1 + F_2 \quad , \quad Q_5 = F_1 - F_2 \tag{13}$$

$$F_i = \int_0^{\infty} dz\, \bar{\chi}_i \gamma_0 \chi_i$$

But this problem - the calculation of fermion number in background field - is well-known [12]. The adiabatic approach [13] to this problem is more useful for our problem. Let us remind, that adiabatic computation gives a reliable prediction only for the fractional part of the fermion number [8, 13]. So, if the background is zero at $t \to -\infty$ $(F_i(t=-\infty)=0)$ and the background have the soliton-profile at $t \to +\infty$, then $\Delta F = \lim_{T\to\infty}[F(T)-F(-T)]$ can be calculated in adiabatic approach and ΔF is equal to the soliton fermion number.

We have interested in variation of $F_{1,2}$ for 2d theory (6) when we pass from one vacuum ($t = -\infty$) to another one ($t = +\infty$). The analogous model have been discussed in ref. [13] and the result is:

$$\Delta F_1 = F_1(T) - F_1(-T) = \pm \frac{1}{2\pi}\left\{ \operatorname{arctg}\frac{\varphi_2}{\varphi_1}(T) - \operatorname{arctg}\frac{\varphi_2}{\varphi_1}(-T)\right\}\Bigg|_{z=0}^{z=\infty} =$$

$$= \pm \frac{1}{2\pi}\left(-\frac{\pi}{2}-\frac{\pi}{2}\right) = \pm\frac{1}{2}$$

Here sign is not determined. It is due to double-degeneration of states. In the language of [13] it corresponds to that we have two solitons with $F_1 = \pm\frac{1}{2}$. We can calculate the value $\Delta F_2 = -\Delta F_1$ analogously. The minus sign is connected with opposite sign of interaction of χ_1, χ_2 with background fields φ_i, A_μ.

By eq. (12) we have:

$$\Delta F = \Delta F_1 + \Delta F_2 = 0 \quad , \quad \Delta Q_5 = \Delta F_1 - \Delta F_2 = \pm 1 \tag{14}$$

Hence the result (11) is derived independently.

Above we have discussed the meron fluctuation connecting two vacua. Let us discuss now another problem: what is the value of axial number of Gribov state irrespective of existence of meron. Remind some properties of this state:

$$\partial_i A_i^a = 0 \quad , \quad A_i^a \frac{\sigma^a}{2} = -i U^+ \partial_i U, \quad U = \exp\left\{ i g(r) \frac{\vec{n}\vec{\sigma}}{2} \right\} \tag{15}$$

$$g(r=0) = 0 \quad , \quad g(r=\infty) = \pm\pi$$

From (13), (15), [13] we have:

$$F_1 = \frac{1}{2\pi} \operatorname{arctg}(-\operatorname{ctg} g(r))\Big|_{r=0}^{r=\infty} = \frac{1}{2\pi}[g(\infty) - g(0)] = \pm\frac{1}{2} \tag{16}$$

$$Q_5 = F_1 - F_2 = \pm 1, \qquad F = F_1 + F_2 = 0$$

Here the values $F_{1,2}$ are determined by function "g" (15). In this calculation both signs occur. In the language of mechanical model [5], fig. 2) it corresponds to both possibilities of motion: $g(\tau=\infty) = \pm\pi$ at $\tau = \ln r = +\infty$, while at $\tau = \ln r = -\infty$, $g(0)=0$, i.e. the particle can move to any minimum of potential.

The result (16) indicates that the normalizable zero-energy eigenmode in the A_i - field (15) exists. This fact can be checked manifestly.

I would like to make some remarks concerning the eqs. (14), (16). In our approach all the vacuum states have $F = F_1 + F_2 = 0$. So, spontaneous breaking of fermion number can not appear. This statement is in agreement with rigorous result of ref. [15] on impossibility of such violations in vector-like theories.

Let us remind that in the fermion number fractionization phenomenon [8,12] the ground state is doubly degenerate and connected with each other by the creating (a^+) and annihilating (a) operators fermion zero energy mode. In our case we have analogous connection:

$$|Q_5 = +1\rangle = a_1^+ a_2 |Q_5 = -1\rangle, \quad |Q_5 = -1\rangle = a_2^+ a_1 |Q_5 = +1\rangle$$
$$\hat{Q}_5 = \left(a_1^+ a_1 - \frac{1}{2}\right) - \left(a_2^+ a_2 - \frac{1}{2}\right) \tag{17}$$

Here a_1^+ is the operator creating zero energy mode of $\chi_1(\psi_L)$ - - field; a_2 is the operator annihilating that of $\chi_2(\psi_R)$ field (5). The chiral number of operator $a_1^+ a_2$ is equal to 2 and fermion number is zero.

Let us note that there is analogous structure of vacuum in the supersymmetric SU(N) gluodinamic [16]. Namely, in this theory we have N vacuum states of spontaneously broken chiral symmetry connected with each other by the discrete chiral rotations: $\psi \to \exp\{i\gamma_5 \pi k/N\}\psi$.

4. Let us derive the eq. (16) without invoking 2d theory (6). The problem can be formulated as follows. We have interaction of Dirac fermions with external static field $A_i = -iU^+\partial_i U$:

$$\mathcal{L}_{int.} = \bar\psi_R \gamma_i (U^+ i\partial_i U)\psi_R + \bar\psi_L \gamma_i (U^+ i\partial_i U)\psi_L . \qquad (18)$$

What is the value of axial number of ground state of this system? In formal mathematical sense analogous problem was solved some years ago. Namely, the problem of calculation of baryon number in the skyrmeon model [17-19] remind formally the problem (18). The differences between these problems are the following ones. a) Instead of baryon number B we calculate axial number Q_5 ; b) instead of π-meson background $V = \exp\{i\frac{\vec\tau}{2}\vec\pi(x)\gamma_5\}$, interacting with ψ_L , ψ_R with opposite signs we have gluon background $U = \exp\{ig(r)\frac{\vec n\vec\sigma}{2}\}$, interacting equally with ψ_L , ψ_R . In the standart adiabatic approach [8,13] one may calculate Q_5 and transform to the form:

$$Q_5 = \frac{1}{24\pi^2}\int d^3x\, \epsilon_{ijk}\, Tr\{(\Omega^+\partial_i\Omega)(\Omega^+\partial_j\Omega)(\Omega^+\partial_k\Omega)\} \qquad (19)$$

$$\Omega = U^2 = \exp\{ig(r)\vec n\vec\sigma\}.$$

The right hand side of eq. (19) can be recognized as the properly normalized integral expression for Pontryagin number in $\pi_3(S^3) = Z$. In a Gribov field (15) the right hand side of (19) equals ± 1 because Ω tends to $\exp\{i\pi\vec n\vec\sigma\}$ when r tends to ∞. It is in agreement with eqs (11,14,16).

I would like to make some remarks concerning the analogy between skyrmeon and Gribov field (15). Remind, that Coulomb gauge for the field (15) consists in the following: $\partial_i(U^+\partial_i U) = 0$.

But it is extremum of action $\int d^4x\, Tr(\partial_\mu U^+)(\partial_\mu U)$ for the static solution of effective chiral theory. Hence, the existence of skyrmenon-soliton and Gribov's solution are equivalent problems. Then, inequality $B \neq 0$ for the skyrmenon is equivalent to statement that $Q_5 \neq 0$ for Gribov's state. The existence of baryon (B = 1) and antybaryon (B = -1) is equivalent to that Gribov's states are doubly degenerate ($Q_5 = \pm 1$).

5. The main result is the following one. The well separated meron-antimeron (topological charge is $\pm$ 1/2) fluctuation with finite action create the Gribov state. This state is doubly degenerate and discrete (S $|Q_5 = +1\rangle \longleftrightarrow |Q_5 = -1\rangle$, $\psi \to e^{i\gamma_5 \frac{\pi}{2}}\psi$ is spontaneously broken.[x] There are the analogous phenomena in 2d theories [6,11], (12) and in the supersymmetric SU (N) gluodynamic [16].

As mentioned above we interpret Gribov's state as soliton, having the normalizable zero energy mode, and connecting two vacua.

Let us note that this phenomenon is possible because transformation $U = \exp\{ig(r)\frac{\vec n\vec\sigma}{2}\}$, $g(0) = 0$, $g(\infty) = \pm\pi$ is unacceptable as gauge one and $-iU^+\partial_i U$ is non-vacuum solution in the presence of fermions. In other words, if ψ - the solution with boundary condition $\psi(x) = \psi(x+L)$, then $U\psi$ is not. It resembles the result of ref. [20] on impossibility to add fermions to the theory with twisted ($U(\infty) \neq U(0)$) boundary conditions. It is due to that the gluon-fields are invariant under the center Z (N) of the gauge group SU(N), but quarks are not invariant under Z (N).

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[x] Let us note that these doubly degenerate states are needed to resolve Θ - puzzle [21]

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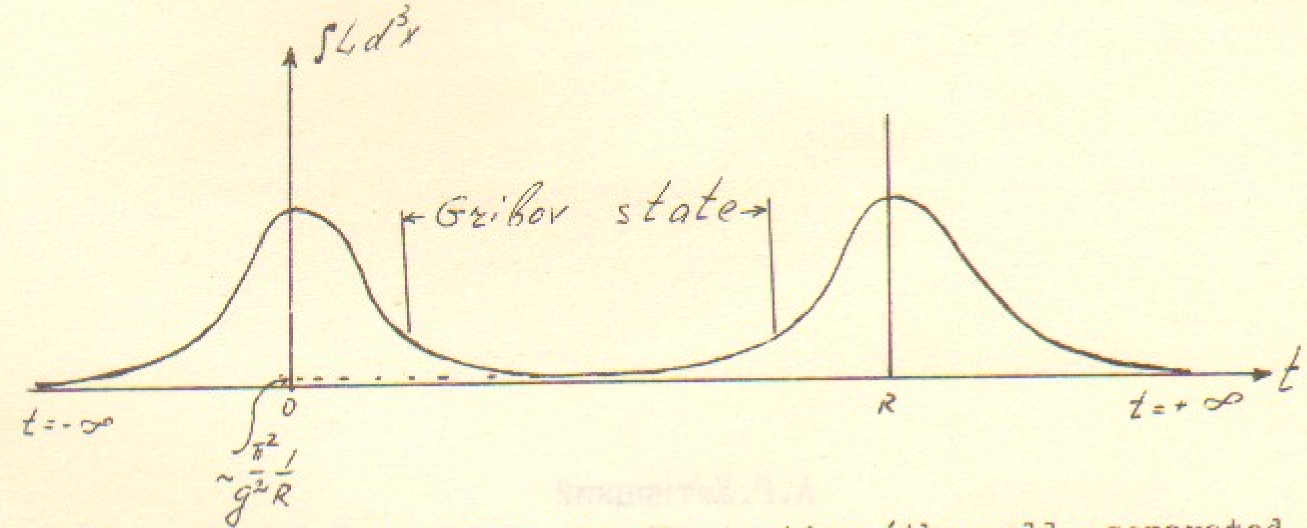

Fig.1 The example of vacuum fluctuation (the well -separated ($R \gg 1$) meron-antimeron) with finite action, which creates quasi-Gribov state.

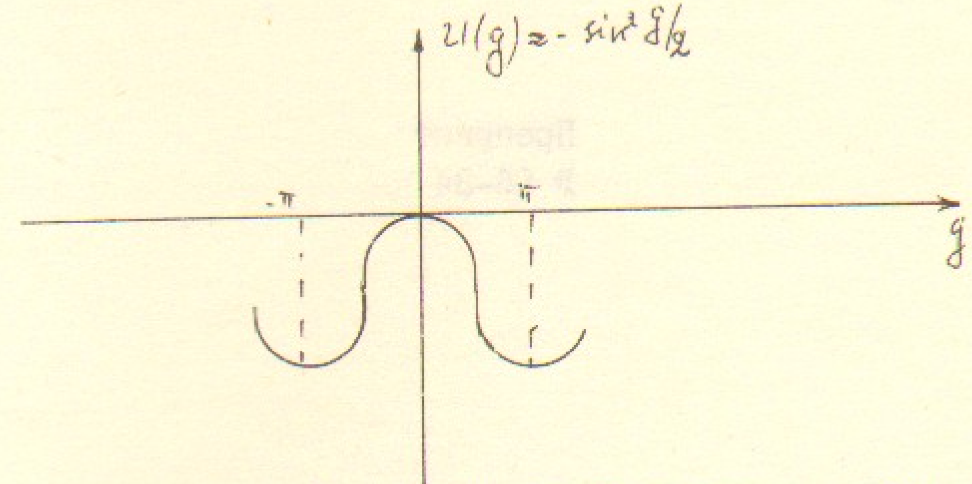

Fig.2 The mechanical model for Gribov equation (15), describing the motion of particle in the potential $U(g) \sim -\sin^2 g/2$

А.Р.Житницкий

НАРУШЕНИЕ ДИСКРЕТНОЙ КИРАЛЬНОЙ СИММЕТРИИ
В КХД КАК ПРОЯВЛЕНИЕ ЭФФЕКТА КАЛЛАНА-РУБАКОВА

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