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HADRON QCD
(BOUND STATES IN GAUGE THEORIES)

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1. Introduction.

Bound states have played a fundamental role in the development of quantum theory. The description of the atomic spectrum by E.Schrödinger signified formation of the quantum mechanics as a consistent theory, and the description of the " the Lamb shift of spectral lines " by H.Bethe became the beginning of the creation of recent QED and quantum field theory [1]

However, there is a belief that the consistent theory of bound states is not constructed up to now. Another belief is that such a theory is not needed in connection with the succesful development of nonperturbative methods of the lattice calculations.

The aim of this paper is to discuss the recent status of bound states in gauge theories and to try to find the additional theoretical and empirical principles of the nonlocal of description of hadrons and atoms.

2. The statement of problem.

We begin with the very known example of an atom in the rest frame with the momentum $\mathcal{P}_\mu = (M_A, 0, 0, 0)$. In the lowest order in the radiative corrections the atom spectrum is described by the action

$$\begin{aligned} W &= \int d^4x \bar{\psi}(x) \left(i\not{\partial} - m^0\right) \psi(x) - \\ &- \frac{1}{2}\int d^4x d^4y \psi(y)\, \bar{\psi}(x)\, \mathcal{K}(x,y)\, \psi(x)\bar{\psi}(y) \end{aligned} \tag{1}$$

where $\mathcal{K}$ is the Coulomb kernel

$$\mathcal{K}(x,y) = (\gamma_0)\cdot(\gamma_0)\, V_C(z)\, \delta(z_0) \;\; ;$$

$z_\mu = (x_\mu - y_\mu)$ is relative space - time. The action (1) leads to the Salpeter equation and eventually to the Schrödinger equation for an atom wave function [2] $\chi(z)$.

[1]The first paper on " the Lamb shift theory " with the result, differing from the H.Bethe formula by factor ~ 1.3, has been reported by D.I.Blokhintsev on the Lebedev Physical Institute seminar in 1938, ten year before the experimental discovery of the Lamb shift. Unfortunally this paper has not been understood and published [1].

The wave function $\chi(z)$ can be used for the construction of the bilocal atom field

$$\Phi(x,y) = \Phi(z|X) = e^{iM_A X_0} \cdot \chi(z) \cdot \delta(z_0) , \quad (2)$$

that depends on two coordinates : the relative $(z_\mu = x_\mu - y_\mu)$ and total $(X_\mu = (x_\mu + y_\mu)/2)$.

An important property of this field is the simultaneity of the elementary particles formed an atom. (A proton yesterday and electron today do not make an atom [3].)

The radiation corrections, breaking the potential simultaneity, do not break the atom bilocal field simultaneity, as it has been shown in ref. [4] .

The main question in the statement of the problem discussed here is "What is the action (1) that describes a moving atom?". The wave function of a relativistic atom (used for the description of the creation of atoms and of their break - down [5]) is constructed by the usual boost operation

$$\Phi(x,y) \longrightarrow \Phi'(x,y) = e^{i\mathcal{P}' \cdot X} \cdot \chi(z^\perp) \cdot \delta(z \cdot \eta') \; ; \; (z^\perp_\mu = z_\mu - \eta'_\mu(z \cdot \eta')) , \quad (3)$$

where $\mathcal{P}'_\mu$ is the total momentum $\mathcal{P}'_\mu = (\sqrt{\mathbf{P}^2 + M_A^2}, \mathbf{P} \neq 0) = M_A \cdot \eta'_\mu$.

This relativistic atom bilocal field is described by the action (1) with the moving Coulomb kernel

$$\mathcal{K}(z|X) = \not\eta' \; V(z^\perp) \; \not\eta' \; \delta(z \cdot \eta') . \quad (4)$$

This means that we choose the new radiative gauge depending on the arbitrary unit time - like vector η'_μ (that one calls by the time - axis of quantization) and this vector has been chosen parallel to the total momentum of an atom ($\eta' \sim \mathcal{P}'_\mu$). According to the Love theorem [4] the old structure of the bilocal field (2) cannot be restored by the radiation corrections. We cannot say here that the atom wave function does not depend on gauge.

So, the usual boost of the matrix elements with the atom wave function corresponds to the Lorentz transformation of field operators accompanied by the gauge change. Just this field transformation law has first been pointed out by Heisenberg and Pauli in 1930 [6]. Another question is " What is the relativistic atom for which we do not change gauge ? "

The answer to this question is given by the general theory of bilocal fields [7,8].

It is easy to see that the bilocal fields (2) and (3) satisfy the Yukawa condition [7]

$$z_\mu \frac{\partial}{\partial X_\mu} \Phi(z|X) = 0 \quad (5)$$

which means that the bilocal field is an irreducible representation of the Lorentz group (i.e. it has the mass $\mathcal{P}^2 = M^2$ and a spin). Expression (5) is a generalized condition of irreducibility of vector, tensor and other fields ($\partial_\mu A_\mu = 0; \partial_\mu T_{\mu\nu} = 0; ...$)[8].

If we shall not change gauge, then the irreducibility condition is not fulfilled and the relativistic dispersion law breaks down, $\mathcal{P}^2 \neq M_A^2$ (see for example ref. [9]). Thus, not only the wave function but also the atom spectrum depend on gauge.

There are several papers[4,10] devoted to the proof of gauge independence of an atom spectrum. In these treatments, the Coulomb interaction is used in the rest frame with the choice of the time-axis $\eta_\mu = (1,0,0,0)$. However, all the authors have not taken into account that the vector η_μ (contained in the Coulomb part of interaction) can indeed be arbitrary, and that a transition from one vector η_μ to another η'_μ $({\eta'_\mu}^2 = 1)$ is realized by means of a special change of the gauge.

One of the reason of the atom physics dependence on gauge consists in that the elementary particles in an atom are off mass - shell.

For example, it is easy to see that the sum of the Coulomb field and transversal photon propagators

$$\mathcal{K}^R(J) = J_0^{(1)} \frac{1}{\mathbf{q}^2} J_0^{(2)} + J_i^{(1)} (\delta_{ij} - q_i \frac{1}{\mathbf{q}^2} q_j) \frac{1}{q_0^2 - \mathbf{q}^2} J_j^{(2)}$$

coincides with the Feynman gauge propagator $\mathcal{K}^F$ up to the longitudinal term $\mathcal{K}^L$:

$$\begin{aligned}
\mathcal{K}^R(J) &\equiv \mathcal{K}^F(J) + \mathcal{K}^L(J), \\
\mathcal{K}^F(J) &= -[J_0^{(1)} J_0^{(2)} - J_i^{(1)} J_i^{(2)}] \frac{1}{q_0^2 - \mathbf{q}^2} \\
\mathcal{K}^L(J) &= ((q_0 J_0^{(1)})(q_0 J_0^{(2)}) - (q_i J_i^{(1)})(q_j J_j^{(2)})) \frac{1}{\mathbf{q}^2(q_0^2 - \mathbf{q}^2)}
\end{aligned}$$

$\mathcal{K}^L$ disappears only on the mass-shell (because of the current conservation law $J_0^{(1,2)} q_0 = J_i^{(1,2)} q_i$). But off the mass-shell for the Bethe-Salpeter equation the currents (J) turn into the vertices (Γ) which do not satisfy the conservation law

$$\Gamma_0^{(1,2)} q_0 \neq \Gamma_i^{(1,2)} q_i ; \quad \mathcal{K}^R(\Gamma) \neq \mathcal{K}^F(\Gamma).$$

We cannot use for the atom description any gauge and any η_μ. From this point of view it is doubtful whether the lattice calculations can describe the Lamb shift.

We have seen that the real action for the relativistic QED atom in the lowest order in coupling constant is given by the action (1) with the kernel (4) where the time - axis η_μ is the unit eigenvector of the bound state total momentum operator (5)

$$\eta'_\mu \Phi(z|X) \sim \frac{\partial}{\partial X_\mu}\Phi(z|X) . \qquad (6)$$

It is wonderful that the relativistic potential model (1), (4), (6) has not been considered until now. One of the problems is the inclusion of a rising potential in this model instead of the Schrödinger equation. Another problem is the foundation of the radiative gauge for the description of the bound state spectrum in the rest frame and the Heisenberg-Pauli group transformation. Just these problems will be discussed below.

3. New relativistic potential model (N.R.P.M.).

N.R.P.M. (1), (4), (6) has been considered in paper [11] by the author with collaborators. This model gives consistent description of massless quarks interacting by means of a rising potential, i.e. it describes the constituent quark mass, spontaneous chiral breaking symmetry, and the massless pion as the Goldstone bilocal mode ($\sqrt{\mathcal{P}^2} = 0$). In the rest frame (for total momentum $\mathcal{P}^2 \neq 0$) our equations for the spectra of quarks and mesons turn into the equations of the model of ref. [9] where the meson spectrum has been got in agreement with the experimental data. We have shown that increasing quark current masses lead to the Schrödinger equation for the heavy quarkonia [11]. Besides the light and heavy quarkonium spectroscopy the N.R.P.M. describes also the hadron interactions.

The well known in the nuclear physics the separable approximation for N.R.P.M. leads to the Nambu - Jona - Lasinio (NJL) model with the definite form factor of the regularization [12]. Thus N.R.P.M. contains also the chiral Lagrangian inspired by one of the versions of the NJL model [13].

4.Gauge dependence of the bound state physics.

First let us recall such notions as "gauge invariance","choice of gauge", and "change of the gauge". The gauge invariance of Lagrangian $\mathcal{L}(A)$ means that it does not vary under gauge transformations

of the fields A

$$\mathcal{L}(A^g) = \mathcal{L}(A) \ , \qquad (7)$$

where

$$A^g = g(A + \partial)g^{-1} \ .$$

The choice of the gauge is a specific gauge transformation g^f depending on the field A , so that the new field $A^f[A] = g^f[A](A+\partial)[g^f[A]]^{-1}$ satisfies the additional condition

$$f(A^f[A]) = 0 \ . \qquad (8)$$

The quantization of the fields and the Feynman rules are always formulated in terms of a certain gauge: $f_1 = 0, f_2 = 0, \dots$. We would like to draw your attention to some not well known consequences of these definitions.

i).The explicit solution of gauge condition (8) gives the physical variables A^f as a functional on the initial fields A_i[6]. In QED this is the axial field

$$A^{(3)}_\mu[A] = (\delta_{\mu\nu} - \partial_\mu \frac{1}{\partial_3}\delta_{3\nu})A_\nu, \quad (A^{(3)}_3[A] = 0),$$

or the transversal field

$$A^T_i[A] = (\delta_{ij} - \partial_i \frac{1}{\partial^2}\partial_j)A_j, \quad (\partial_i A^T[A] = 0)$$

and so on.These functionals are invariant under gauge transformations of the initial fields in the sense of eq.(4). So, *any gauge choice is a transition from the initial fields to the gauge invariant physical variables* (i.e. " gauge " (8) is the choice of variables).

ii).The change of the gauge (from A^{f_1} to A^{f_2}) is fulfilled by the substitution [7]

$$\begin{aligned} A^{f_2}[A^{f_1}] &= V[A^{f_1}](A^{f_1} + \partial)V^{-1}[A^{f_1}]; \\ \psi^{f_2} &= V[A^{f_1}]\psi^{f_1} \end{aligned} \qquad (9)$$

All Green functions are invariant under the operation (9)

$$< \psi^{f_2} \dots \bar{\psi}^{f_2} > \equiv < V[A^{f_1}]\psi^{f_1} \dots \bar{\psi}^{f_1} V^{-1}[A^{f_1}] > \qquad (10)$$

(if anomalies are absent). This substitution contains not only the modification of the Feynman rules (i.e.the gauge change) but also the spurious diagrams induced by the factor $V[A^{f_1}]$ (which do not follow from the initial Lagrangian).

On the mass-shell these additional diagrams do not contribute, and the invariance under the gauge change takes place. But off the mass- shell the dependence on the gauge takes place and this does not mean the gauge noninvariance (any variables (8) are gauge invariant as we have seen above).

The experimental value of the Lamb - shift is described by any " gauge " up to the spurious diagrams. The radiative " gauge " is unique which does not demand these spurious diagrams to reproduce the observed Lamb - shift in atomic spectra.

5. The minimal quantization scheme.

The Feynman rules in the radiative gauge applied in the atomic physics and the Heisenberg-Pauli relativistic group can be justified by the minimal quantization scheme of gauge field theories which has been formulated in ref.[14] as the following two axioms:

i).The axiom of the choice of physical variables by the projection of the Lagrangian and the Belinfante energy-momentum tensor

$$T_{\mu\nu} = F_{\mu\lambda}F^{\lambda}_{\nu} + \bar{\psi}\gamma_{\mu}[i\partial_{\nu} + eA_{\nu}]\psi - g_{\mu\nu}L + \frac{i}{4}\partial_{\lambda}[\bar{\psi}\Gamma^{\lambda}_{\mu\nu}\psi] , \qquad (11)$$

$$\Gamma^{\lambda}_{\mu\nu} = \frac{1}{2}[\gamma^{\lambda}\gamma_{\mu}]\gamma_{\nu} - g_{\mu\nu}\gamma^{\lambda} - g^{\lambda}_{\nu}\gamma_{\mu}$$

upon the Gauss equation solution for the time component $A_0 = (\eta \cdot A)$

$$\frac{\partial L}{\partial A_0} = 0.$$

ii).The axiom of quantization of the minimal set of physical variables by the diagonalization of the Belinfante Hamiltonian T_{00}.

In QED the first axiom expresses the tensor (11) only in terms of the transversal variables A^T, ψ^T as nonlocal gauge invariant functionals on the initial fields

$$T_{\mu\nu}[A_i, A_0 = (\frac{1}{\partial^2}\partial_i\partial_0 A_i + j_0)] = T_{\mu\nu}[A^T_i[A_i], \psi^T[A,\psi]] . \qquad (12)$$

Here

$$\hat{A}^T_i[A] = V[A](\hat{A}_i + \partial_i)V^{-1}[A] ,$$
$$\psi^T[A,\psi] = V[A]\psi ,$$
$$V[A] = exp(\frac{1}{\partial^2}\partial_i\hat{A}_i) ,$$
$$\hat{A} = ieA .$$

The usual Lorentz transformation of the initial fields in the Gauss equation leads to the Heisenberg--Pauli transformaton of the transversal functional

$$\psi^T[A_i + \delta^0_L\psi, \psi + \delta^0_L\psi] - \psi^T[A_i,\psi] = \delta^0_L\psi^T + ie\Lambda\psi^T , \qquad (13)$$

where ϵ_k are the transformation parameters

$$\delta^0_L = [\epsilon_i(x_i\partial_t - t\partial_i) + \epsilon_k\gamma_0\gamma_k]; \quad \Lambda = \epsilon_k\frac{1}{\partial^2}[\partial_0 A^T_k + \frac{\partial_k}{\partial^2}j_0] . \qquad (14)$$

The second axiom leads to the same transformation law (13) for quantum fields

$$i\epsilon_k[\int dx(T_{00}x_k - T_{0k}t), \psi^T] = \delta^0_L\psi^T + ie\Lambda\psi^T .$$

In the minimal quantization scheme the relativistic transformation of the classical variables (12) coincides with the quantum ones on the operator level.

This coincidence is the main difference between the minimal quantization and the one in the usual radiative gauge.Another difference is the phase physics due to the infrared zero modes in the exponent of the factor $V[A]$ in eq.(12).

The same explicit construction of the physical variables for non-Abelian theory[14,15] leads to the topological degeneration of these phase factors and to a confinement mechanism as a destructive phase interference.

The third difference from the conventional Dirac approach is the dependence of the bound state physics on the time-axis of quantisation η_μ and the importance of one more empirical bound state principle, the Markov-Yukawa choice of the time-axis (6).

The minimal quantization with the Markov-Yukawa choice of the time-axis η_μ does not change the S- matrix with asymptotical free states of elementary particles (as this S- matrix does not depend on gauge and on η_μ) but these empirical axioms are necessary and really are used in the atomic physics independently of the validity of perturbation theory[4].

6. New QCD myth.

The minimal quantization of chromodynamics[14] up to the phase phenomenon [15] is reduced to the explicit gauge invariant construction of the Schwinger operator quantization of the non-Abelian theory[16] with the Hamiltonian

$$\begin{aligned}\mathcal{H}(g^2D) &= \int d\mathbf{x}[\frac{1}{2}(E_i^a(\mathbf{x}))^2 + \frac{1}{4}(F_{ij}^a(\mathbf{x}))^2 + \bar{q}(\mathbf{x})(i\gamma_k\nabla_k + m^0)q(\mathbf{x})] + \\ &+ \frac{1}{2}\int d\mathbf{x}d\mathbf{y}J_{tot}^a(\mathbf{x})[g^2D^{ab}(\mathbf{x}-\mathbf{y}|A)]J_{tot}^b(\mathbf{y}) + \\ &+ \ nonlocal\ Schwinger\ terms\end{aligned} \quad (15)$$

Here

$$\begin{aligned}\nabla_k &= \partial_k + gA_k^a\frac{\lambda^a}{2};\ F_{ij}^a = \partial_iA_j^a - \partial_jA_i^a + gf^{abc}A_i^bA_j^c; \\ J_{tot}^a &= q^+\frac{\lambda^a}{2}q + f^{abc}E_i^bA_i^c;\ \partial_iA_i^a = \partial_iE_i^a = 0\end{aligned} \quad (16)$$

g is the coupling constant and the function $D^{ab}(\mathbf{x}-\mathbf{y}|A)$ satisfies the equation

$$[(\nabla_i\partial_i)\frac{1}{\partial^2}(\nabla_j\partial_j)]^{ac}D^{cb}(\mathbf{x}-\mathbf{y}|A) = \delta(\mathbf{x}-\mathbf{y})\delta^{ab} \quad (17)$$

(Where $\nabla_i^{ab} = \delta^{ab}\partial_i + gf^{abc}A_i^c$). The Schwinger terms are defined from the Lorentz covariance condition [14,16].

We shall consider the Hamiltonian (15) as a basis for construction of QCD for hadrons.

Just this Hamiltonian (unlike the QED one) contains a new type of infrared divergences at zero three-dimensional momenta $\mathbf{k}^2 = 0$.

The asymptotic freedom formula cannot remove these static divergences and becomes a phenomenological supposition. The removal of these divergences has not only a purely mathematical (theoretical) character .(Recall that in QED the solution of the infrared problem is accompanied by including a phenomenological parameter of the type of the dimension of a device.)

One thing is known: these static divergences bear a relation to the modification of the static Coulomb potential at long distances (or at $\mathbf{k}^2 \sim 0$) and to the physical dimensional transmutation. Instead of the asymptotic freedom phenomenology let us take the form and the parameter of the modification from the experiment: i.e. the heavy quarkonium spectroscopy that definitely points out the rising potential[17]. (This potential can be forced by the nontrivial boundary condition of the Gauss equation[15] like Λ_{QCD} appeared in the boundary condition of the renormalization group equations.)

We would like to draw your attention to the wonderful fact: that the rising potential ansatz

$$\mathcal{H}(g^2D(\mathbf{x}|A)) \rightarrow \mathcal{H}(V_R(\mathbf{x}) + g^2D(\mathbf{x}|A)) \quad (18)$$

removes all infrared divergences in a perturbation theory in the coupling constant g^2[18,19].This hadron QCD perturbation theory contains in particular the old parton QCD, the nonlocal chiral Lagrangian for light quarks, and the potential model for J/ψ spectroscopy.

We comment here on some details of the hadron QCD (QCD_h) [19]. We choose as a test potential the oscillator one with the dimension parameter $\sim$ 300 MeV. In the lowest order in coupling constant the rising potential leads to the constituent masses of light quarks and gluons and does not change the heavy quark masses[11,18]. The QCD_h perturbation theory is formulated in terms of the modified gluon and quark propagators which in the explicit form depend on the total hadron momentum $\mathcal{P}_\mu$. For large transverse momenta $|q^\perp| \geq$300 MeV these modified propagators turn into the parton ones of the usual QCD without confinement properties.

The modified gluon propagator also modifies the running coupling constant in the region of small transfer momenta[18]. The new running coupling constant has no singularities in the whole region of transfer momenta and is smaller than $\alpha_s^{mod}(0) \sim 0.2$. At large momenta it coincides with the asymptotic freedom formula. QCD_h describes the glueball masses in the region expected now[18].

7. Phenomenology and theory of confinement.

In $QCD_{(h)}$ we face with the continuous quark and gluon spectra in spite of the rising potential (see also[9,20]). This means that the quark - gluon loops have the imaginary parts which contradict the potential confinement, and vice versa, the absence of the imaginary parts contradicts the phenomenology of the measurement of quark and gluon quantum numbers that is based on the parton interpretation of deep-inelastic process. Its essence consists in that the sum over all hadron final states is described as an imaginary part of the corresponding elastic amplitude constructed from quark-gluon diagrams of the QCD perturbation theory.

This description is called the quark-hadron duality (QHD) and is used in the phenomenology as the energy averaging (the global QHD) and without averaging, in the energy region far from resonances (the local QHD). For example, the cross section of the process e^+e^- into hadrons in the nonresonance energy region not only on the average but also at points coincides with the imaginary part of quark-gluon loops.

The local QHD means that the perturbation theory is realy used in the Minkowski space. To get the QHD, it is sufficient to remove quark and gluon (color) states from physical ones in the unitarity relation for the S-matrix.

In QHD there are implicitly used two different types of quarks states : the physical states (c) for which $T_{ic} = 0$ and parton states (p) $T_{ip} \neq 0$, which reflect only particular analytical properties of "elastic" *hadron* amplitudes reproduced by the imaginary parts of quarks diagrams.

One of the formulations of the confinement problem is as follows : Why doesn't the coincidence of physical and parton states occur in QCD ? Another formulation : Why is the probability of color particle production equal to zero , $T_{ic} = 0$, while the probability of hadronization is equal to unity ?

It is wonderful that the theoretical observation of quarks as partons and their experimental nonobservation take place in the same energy region of the Minkowski space. This is the main paradox of the parton phenomenology and QHD.

The answer to the question : " How can the perturbation theory in the Minkowski space ($T_{ip} \neq 0$) be made consistent with the confinement hypothesis ($T_{ic} = 0$) ? " is given neither by asymptotical freedom, nor by the confinement potential.

The first explains only the QCD perturbation theory in the Euclidean space (where theoretical quantities are connected with the realistic cross sections by the dispersion relation,or the energy averaging).

The argumentation of the potential version of confinement is based only on different regimes of the quark behaviour in different (not the same) energy regions. Moreover, all attempts to explain the

nonobservability of individual quarks by solving the Dyson-Schwinger equation for quark propagator have led not to confinement, but, rather to the spontaneous breaking of chiral symmetry.

The self-consistent solution of the quark-parton paradox has been given first by the t'Hooft two-dimensional chromodynamics [21] , where the confinement of physical quarks is explained not by the interaction potential,but by the process of dressing bare quarks.In this model [21] all physical quarks have infinite masses as a consequence of the infrared divergences whereas colorless amplitudes are expressed in terms of bare quark propagators with finite masses without infrared divergences.

The absence of the amplitude of color particle production does not contradict the unitarity condition in the Minkowski space.The point is that when bound states are present,the unitarity relation should not be understood as an identity,but rather as one of the self-consistency conditions of the theory used for normalizing the bound states and their interaction constant.If for some reason the color states disappear, the probability of other channels increases so that the total probability is equal to unity.

A similar confinement mechanism of "dressing" bare quarks is contained in the minimal scheme of quantization of the non-Abelian theory[14] due to the topological degeneration of the physical variables [15].

8. Summary.

i).We have distinguished between the free asymptotic state S-matrix (S_F) and the bound asymptotical state S-matrix (S_B). S_B depends on gauge,unlike S_F (we have seen that the gauge change does not coincide with the gauge transformation, and the gauge dependence is not the gauge non-invariance). For the proof of the relativistic covariance of the free asymptotical state S-matrix it is sufficient to pass to any relativistic gauge.For the relativistic description of S_B , *instead of the Dirac quantization* method with the canonical gauge-noninvariant Hamiltonian and an arbitrary gauge, *we use the minimal method* with the choice of the time-axis along the eigenvector of the bound state total momentum operator.

ii). In the low-energy region we give up the asymptotic freedom formula which in this region goes out of the range of validity and turns into the phenomenological supposition. We decline the renormalization group equations as a cause of the infrared dimensional transmutation since these equations bear a relation to the ultra-violet divergences rather than to the infrared ones (recall that in the ultraviolet finite theories the renormalizable group equations turn into the trivial identities and do not contain physical information[21]).

Instead of the renormalization group dimensional transmutation we use the rising potential - ansatz as the infrared physical regularization of the perturbation theory.

iii). We have seen that the *rising potential leads to constituent masses, hadron spectroscopy, chiral Lagrangians but not to confinement.*

Instead of the potential confinement we have the destructive interference phenomenon [15] which is possible in the minimal quantization method and which leads just to the quark - hadron duality formula where the parton states differ from the physical color ones like in the t'Hooft two-dimensional QCD[21].

In conclusion we would like to note two intriguing questions: What is the real physical cause of the rising potential? and What is the time-axis for the Universe as a bound state of all their particles?

It is our belief that the answer to the first question lies beyond QCD and has no relation to the renormalization equation.

As to the second question, if the "time-axis" η_μ is a eigenvalue of the operator of the derivative with respect to the total coordinate of all particles in the Universe including any man, then any motion of the man gives a contribution to the "time".

We finish with the definition of the human freedom belonging to the father of the Byzatine Theology St. Maximus (580-662)[23]:"*Any motion*, if it makes sense, possesses also *a freedom*, and the task of it is the realization of a good moral existence, the final aim of which will be the "sense of an everlasting existence".

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Первушин В.Н. E2-89-604
Адронная КХД /Связанные состояния в калибровочных теориях/

Предлагаются общие принципы описания связанных состояний в КЭД и КХД с целью простроения самосогласованной схемы вычисления спектра и амплитуд взаимодействий адронов. Такими принципами являются явное решение уравнений Гаусса на временную компоненту, квантование минимального набора физических переменных и выбор времени квантования в соответствии с релятивистской теорией билокальных полей Маркова-Юкавы. Построенна по этим принципам КХД содержит новые инфракрасные расходимости, меняющие поведение кулоновского поля на больших расстояниях. Эти расходимости /как и инфракрасные расходимости в КЭД/ устраняются с помощью феноменологии: в данном случае, учетом растущего потенциала как "непертурбативного" фона для новой теории возмущений. Показано, как в такой теории адронов возникают партонная модель, нерелятивистская потенциальная спектроскопия, киральные лагранжианы и конфайнмент. Метод квантования Дирака, ренормгрупповые уравнения и вычисления на решетках в их общепринятой формулировке оказываются несостоятельными для описания связанных состояний.

Работа выполнена в Лаборатории теоретической физики ОИЯИ.

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Pervushin V.N. E2-89-604
Hadron QCD (Bound States in Gauge Theories)

The general principles of the description of bound states in QED and QCD are proposed for the aim of construction of the consistent scheme of calculating hadron spectrum and interaction amplitudes. Such principles are the explicit solution of the Gauss equation for time component, the quantization of the minimal set physical variables and the choice of the time-axis of quantization in accordance with the Markov-Yukawa relativistic theory of bilocal fields. QCD constructed by these principles contains new infrared divergences, changing the behaviour of the Coulomb field on large distances. These diverances (like ones in QED) are removed out with the help of phenomenology, in this case, by taking into account the rising potential as the "non-perturbative background" for a new perturbation theory. It is shown how in such hadron theory the parton model, nonrelativistic potential spectroscopy, chiral Lagrangian and confinement appear. The Dirac quantization method, renormalization group equations and lattice calculations in their conventional formulation are proved to be untenable for the description of bound states.

The investigation has been performed at the Laboratory of Theoretical Physics, JINR.

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