Abstract

"Equivalent unconstrained systems" for QCD obtained by resolving the Gauss law are discussed. We show that the effects of hadronization, confinement, spontaneous chiral symmetry breaking and \( \eta_0 \)-meson mass can be hidden in solutions of the non-Abelian Gauss constraint in the class of functions of topological gauge transformations, in the form of a monopole, a zero mode of the Gauss law, and a rising potential.

(Key-words: QCD, Gauss law, topology, monopole, zero mode, hadronization, confinement, U(1)-problem)

1 Introduction

The consistent dynamic description of gauge constrained systems was one of the most fundamental problems of theoretical physics in the 20th century. There is an opinion that the highest level of solving the problem of quantum description of gauge relativistic constrained systems is the Faddeev-Popov (FP) integral for unitary perturbation theory [2]. In any case, just this FP integral was the basis to prove renormalizability of the unified theory of electroweak interactions in papers by ’t Hooft and Veltman marked by the 1999 Nobel prize.

Another opinion is that the FP integral has only the intuitive status. The most fundamental level of the description of gauge constrained systems is the derivation of "equivalent unconstrained systems" compatible with the simplest variation methods of the Newton mechanics and with the simplest quantization by the Feynman path integral. It was the topic of Faddeev’s paper [1] "Feynman integral for singular Lagrangians" where the non-Abelian "equivalent unconstrained system" was obtained (by explicit resolving the Gauss law), in order to justify the intuitive FP path integral [4] by its equivalence to the Feynman path integral. Faddeev managed to prove the equivalence of the Feynman integral to the FP one only for the scattering amplitudes [1] where all particle-like excitations of the fields are on their mass-shell. However, this equivalence is not proved and becomes doubtful for the cases of bound states, zero modes and other collective phenomena where these fields are
off their mass-shell. It is just the case of QCD. In this case, the FP integral in an arbitrary relativistic gauge can lose most interesting physical phenomena hidden in the explicit solutions of constraints [3, 4].

The present paper is devoted to the derivation an ”equivalent unconstrained systems” for QCD in the class of functions of topologically nontrivial transformations, in order to present here a set of arguments in favor of that physical reasons of hadronization and confinement in QCD can be hidden in the explicit solutions of the non-Abelian constraints.

2 Equivalent Unconstrained Systems in QED

The Gauss law constraint is the equation for the time component of a gauge field

$$\frac{\delta W}{\delta A_0} = 0 \Rightarrow \partial_0^2 A_0 = \partial_k \dot{A}_k + J_0$$

in the frame of reference with an axis of time \(l^{(0)}(0) = (1, 0, 0, 0)\). Heisenberg and Pauli [5] noted that the gauge \(\partial_k A_k^* \equiv 0\) is distinguished, and Dirac [6] constructed the corresponding (”dressed”) variables \(A^*_k\) in the explicit form

$$ie A^*_k = U(A)(ieA_k + \partial_k)U(A)^{-1}, \quad U(A) = \exp[ie\frac{1}{\partial_0^2} \partial_k A_k],$$

using for the phase the time integral of the spatial part of the Gauss law \(\partial_k \dot{A}_k\). The action for an equivalent unconstrained system (EUS) for QED is derived by the substitution of the solution of the Gauss law in terms of the ”dressed” variables into the initial action

$$W_{\text{Gauss–shell}} = W_{l^{(0)}(0)}^*(A^*, E^*) .$$

The peculiarity of the ”equivalent unconstrained system” for QED is the electrostatic phenomena described by the monopole class of functions \((f(x) = O(1/r), |x| = r \to \infty)\).

The ”equivalent unconstrained system” can be quantized by the Feynman path integral in the form

$$Z_F[l^{(0)}, J^*] = \int d^2 A^* d^2 E^* \exp \left\{ iW_{l^{(0)}}^*[A^*, E^*] + i \int d^4 x [J^*_k \cdot A_k^* - J^*_0 \cdot A_0^*] \right\}$$

where \(J^*\) are physical sources. This path integral depends on the axis of time \(l^{(0)}(0) = (1, 0, 0, 0)\).

One supposes that the dependence on the frame \((l^{(0)})\) can be removed by the transition from the Feynman integral of ”EUS” (3) to perturbation theory in any relativistic-invariant gauge \(f(A) = 0\) with the FP determinant

$$Z_{FP}[J] = \int d^4 A \delta[f(A)] \Delta_{FP} \exp \left\{ iW[A] - i \int d^4 x J^*_0 \cdot A^* \right\}.$$  

This transition is well-known as a ”change of gauge”, and it is fulfilled in two steps

I) the change of the variables \(A^*\) (4), and

II) the change of the physical sources \(J^*\) of the type of

$$A_k^*(A) J_k^* = U(A) \left( A_k - \frac{i}{e} \partial_k \right) U^{-1}(A) J_k^* \Rightarrow A_k J^k .$$
At the first step, all electrostatic monopole physical phenomena are concentrated in the Dirac gauge factor \( U(A) \) that accompanies the physical sources \( J^* \).

At the second step, changing the sources \( \delta \) we lose the Dirac factor together with the whole class of electrostatic phenomena including the Coulomb-like instantaneous bound state formed by the electrostatic interaction.

Really, the FP perturbation theory in the relativistic gauge contains only photon propagators with the light-cone singularities forming the Wick-Cutkosky bound states with the spectrum differing \( \delta \) from the observed one which corresponds to the instantaneous Coulomb interaction. Thus, the restoration of the explicit relativistic form of EUS \( l^{(0)} \) by the transition to a relativistic gauge loses all electrostatic "monopole physics" with the Coulomb bound states.

In fact, a moving relativistic atom in QED is described by the usual boost procedure for the wave function, which corresponds to a change of the time axis \( l^{(0)} \to l \), i.e., motion of the Coulomb potential \( W_C \) itself

\[
W_C = \int d^4x d^4y \frac{1}{2} J_t(x) V_C(z^\perp) J_t(y) \delta(l \cdot z) ,
\]

where \( J_t = l_\mu J^\mu \), \( z^\perp = z_\mu - l_\mu (z \cdot l) \), \( z_\mu = (x - y)_\mu \). This transformation law and the relativistic covariance of EUS in QED has been predicted by von Neumann \( \delta \) and proven by Zumino \( \delta \) on the level of the algebra of generators of the Poincaré group. Thus, on the level of EUS, the choice of a gauge is reduced to the choice of a time axis (i.e., the reference frame). A time axis is chosen to be parallel to the total momentum of a bound state, so that the coordinate of the potential always coincides with the space of the relative coordinates of the bound state wave function to satisfy the Yukawa-Markov principle \( \delta \) and the Eddington concept of simultaneity ("yesterday's electron and today's proton do not make an atom") \( \delta \).

In other words, each instantaneous bound state in QED has a proper EUS, and the relativistic generalization of the potential model is not only the change of the form of the potential, but sooner the change of a direction of its motion in four-dimensional space to lie along the total momentum of the bound state. The relativistic covariant unitary perturbation theory in terms of instantaneous bound states has been constructed in \( \delta \).

### 3 Unconstrained QCD

#### 3.1 Topological degeneration and class of functions

We consider the non-Abelian \( SU_c(3) \) theory with the action functional

\[
W = \int d^4x \left\{ \frac{1}{2}(G_{0i}^a)^2 - B_i^{a2}) + \bar{\psi}[i\gamma^\mu(\partial_\mu + \hat{A}_\mu) - m]\psi \right\} ,
\]

where \( \psi \) and \( \bar{\psi} \) are the fermionic quark fields. We use the conventional notation for the non-Abelian electric and magnetic fields

\[
G_{0i}^a = \partial_0 A_i^a - D_i^{ab}(A) A_b^a , \quad B_i^a = \epsilon_{ijk} \left( \partial_j A_k^a + \frac{g}{2} f^{abc} A_j^b A_k^c \right) ,
\]

\( ^1 \)The author thanks W. Kummer who pointed out that in Ref. \( \delta \) the difference between the Coulomb atom and the Wick-Cutkosky bound states in QED has been demonstrated.
as well as the covariant derivative $D_i^{ab}(A) := \delta^{ab}\partial_i + gf^{acb}A_c^i$.

The action (8) is invariant with respect to gauge transformations $u(t, \vec{x})$

$$\hat{A}_i^u := u(t, \vec{x}) \left( \hat{A}_i + \partial_i \right) u^{-1}(t, \vec{x}), \quad \psi^u := u(t, \vec{x})\psi ,$$

where $\hat{A}_\mu = g^{2\pi A}_\mu$.

It is well-known [12] that the initial data of all fields are degenerated with respect to the stationary gauge transformations $u(t, \vec{x}) = v(\vec{x})$. The group of these transformations represents the group of three-dimensional paths lying on the three-dimensional space of the $SU(3)$-manifold with the homotopy group $\pi_{(3)}(SU(3)) = \mathbb{Z}$. The whole group of stationary gauge transformations is split into topological classes marked by the integer number $n$ (the degree of the map) which counts how many times a three-dimensional path turns around the $SU(3)$-manifold when the coordinate $x_i$ runs over the space where it is defined. The stationary transformations $v^n(\vec{x})$ with $n = 0$ are called the small ones; and those with $n \neq 0$

$$\hat{A}_i^{(n)} := v^n(\vec{x})\hat{A}_i(\vec{x})v^n(\vec{x})^{-1} + L_i^n , \quad L_i^n = v^n(\vec{x})\partial_i v^n(\vec{x})^{-1} ,$$

the large ones.

The degree of a map

$$\mathcal{N}[n] = -\frac{1}{24\pi^2} \int d^3x \epsilon^{ijk} Tr[L_i^n L_j^n L_k^n] = n .$$

as the condition for normalization means that the large transformations are given in the class of functions with the spatial asymptotics $O(1/r)$. Such a function $L_i^n$ ([1]) is given by

$$v^n(\vec{x}) = \exp(n\hat{\Phi}_0(\vec{x})), \quad \hat{\Phi}_0 = -i\pi\frac{\lambda^a}{r}f_0(r) ,$$

where the antisymmetric SU(3) matrices are denoted by

$$\lambda^1_A := \lambda^2, \quad \lambda^2_A := \lambda^5, \quad \lambda^3_A := \lambda^7 ,$$

and $r = |\vec{x}|$. The function $f_0(r)$ satisfies the boundary conditions

$$f_0(0) = 0, \quad f_0(\infty) = 1 ,$$

so that the functions $L_i^n$ disappear at spatial infinity $\sim O(1/r)$. The functions $L_i^n$ belong to monopole-type class of functions. It is evident that the transformed physical fields $\hat{A}_i$ in (11) should be given in the same class of functions.

The statement of the problem is to construct an equivalent unconstrained system (EUS) for the non-Abelian fields in this monopole-type class of functions.

### 3.2 The Gauss Law Constraint

So, to construct EUS, one should solve the non-Abelian Gauss law constraint [3, 13]

$$\frac{\delta W}{\delta A_0^c} = 0 \Rightarrow (D^2(A))^{ac}A_0^c = D_i^{ac}(A)\partial_0 A_i^c + j^a_0 ,$$

where $j^a_\mu = g\bar{\psi} \frac{\gamma^a}{2} \gamma_\mu \psi$ is the quark current.
As dynamical gluon fields $A_i$ belong to a class of monopole-type functions, we restrict ourselves to ordinary perturbation theory around a static monopole $\Phi_i(\vec{x})$

$$A_i^c(t, \vec{x}) = \Phi_i^c(\vec{x}) + \bar{A}_i^c(t, \vec{x}) ,$$

where $\bar{A}_i$ is a weak perturbative part with the asymptotics at the spatial infinity

$$\Phi_i(\vec{x}) = O\left(\frac{1}{r}\right), \quad \bar{A}_i(t, \vec{x})|_{\text{asymptotics}} = O\left(\frac{1}{r^{1+l}}\right) \quad (l > 1) .$$

We use, as an example, the Wu-Yang monopole $[14, 15]$

$$\Phi_{WY}^i = -i\frac{\lambda^a}{2} \epsilon_{ia}^k \frac{x^k}{r^2} f_{WY}^1 , \quad f_{WY}^1 = 1$$

which is a solution of classical equations everywhere besides the origin of coordinates. To remove a singularity at the origin of coordinates and regularize its energy, the Wu-Yang monopole is changed by the Bogomol’nyi-Prasad-Sommerfield (BPS) monopole $[16]$

$$f_{WY}^1 \Rightarrow f_{BPS}^1 = \left[1 - \frac{r}{\epsilon \sinh(r/\epsilon)}\right], \quad \int d^3x [B^a_i(\Phi_k)]^2 = \frac{4\pi}{g^2\epsilon} ,$$

(19)

to take the limit of zero size $\epsilon \to 0$ at the end of the calculation of spectra and matrix elements. This case gives us the possibility to obtain the phase of the topological transformations (13) in the form of the zero mode of the covariant Laplace operator in the monopole field

$$(D^2)^{ab}(\Phi_{BPS}^A)(\Phi_{BPS}^0)^b(\vec{x}) = 0 .$$

(20)

The nontrivial solution of this equation is well-known $[16]$; it is given by equation (13) where

$$f_{BPS}^0 = \left[\frac{1}{\tanh(r/\epsilon)} - \frac{\epsilon}{r}\right] ,$$

(21)

with the boundary conditions (14). This zero mode signals about a topological excitation of the gluon system as a whole in the form of the solution $Z^a$ of the homogeneous equation

$$(D^2(A))^a_b Z^b = 0 ,$$

(22)

i.e., a zero mode of the Gauss law constraint $[15] [13] [17]$ with the asymptotics at the space infinity

$$\hat{Z}(t, \vec{x})|_{\text{asymptotics}} = \hat{N}(t)\hat{\Phi}_0(\vec{x}) ,$$

(23)

where $\hat{N}(t)$ is the global variable of this topological excitation of the gluon system as a whole. From the mathematical point of view, this means that the general solution of the inhomogeneous equation $[13]$ for the time-like component $A_0$ is a sum of the homogeneous equation (22) and a particular solution $\bar{A}_0^a$ of the inhomogeneous one (13): $A_0^a = Z^a + \bar{A}_0^a$.

The zero mode of the Gauss constraint and the topological variable $\hat{N}(t)$ allow us to remove the topological degeneration of all fields by the non-Abelian generalization of the Dirac dressed variables $[2]$

$$0 = U_Z(\hat{Z} + \partial_0)U_Z^{-1} , \quad \hat{A}_i^a = U_Z(\hat{A}_i^T + \partial_t)U_Z^{-1} , \quad \psi^a = U_Z\psi^T ,$$

(24)
where the spatial asymptotics of $U_Z$ is
\[ U_Z = T \exp\left[ \int dt' \tilde{Z}(t', \vec{x}) \right] \text{asymptotics} = \exp[N(t) \tilde{\Phi}_0(\vec{x})] = U_{as}^{(N)} , \] (25)
and $A^I = \Phi + \bar{A}, \psi^I$ are the degeneration free variables with the Coulomb-type gauge in the monopole field
\[ D_{ac}^{\alpha}(\Phi) \bar{A}_c^\alpha = 0 . \] (26)
In this case, the topological degeneration of all color fields converts into the degeneration of only one global topological variable $N(t)$ with respect to a shift of this variable on integers: $(N \Rightarrow N + n, n = \pm 1, \pm 2, ... )$. One can check \[18\] that the Pontryagin index for the Dirac variables (24) with the asymptotics (17), (23), (25) is determined by only the difference of the final and initial values of the topological variable
\[ \nu[A^*] = \frac{g^2}{16\pi^2} \int_{t_{in}}^{t_{out}} dt \int d^3x G_{\mu\nu}^a G^{a\mu\nu} = N(t_{out}) - N(t_{in}) \] (27)
The considered case corresponds to the factorization of the phase factors of the topological degeneration, so that the physical consequences of the degeneration with respect to the topological nontrivial initial data are determined by the gauge of the sources of the Dirac dressed fields $A^*, \psi$
\[ W_{l(0)}^*(A^*) = \int d^4x J^c^* A^c^* = W_{l(0)}^*(A^I) + \int d^4x J^c^* A^c^*(A^I) . \] (28)
The nonperturbative phase factors of the topological degeneration can lead to a complete destructive interference of color amplitudes \[3, 17, 19\] due to averaging over all parameters of the degenerations, in particular
\[ <1|\psi^*|0> = <1|\psi^I|0> \lim_{L \to \infty} \frac{1}{2L} \sum_{n=-L}^{n=L} U_{as}^{(n)}(x) = 0 . \] (29)
This mechanism of confinement due to the interference of phase factors (revealed by the explicit resolving the Gauss law constraint \[3\]) disappears after the change of “physical” sources $A^*J^* \Rightarrow AJ$ (that is called the transition to another gauge). Another gauge of the sources loses the phenomenon of confinement, like a relativistic gauge of sources in QED \[6\] loses the phenomenon of electrostatics in QED.

### 3.3 Physical Consequences

The dynamics of physical variables including the topological one is determined by the constraint-shell action of an equivalent unconstrained system (EUS) as a sum of the zero mode part, and the monopole and perturbative ones
\[ W_{l(0)} = W_{\text{Gauss-shell}} = W_Z[N] + W_{\text{mon}}[\Phi_i] + W_{\text{loc}}[\bar{A}] . \] (30)
The action for an equivalent unconstrained system \[30\] in the gauge \[26\] with a monopole and a zero mode has been obtained in the paper \[18\] following the paper \[1\]. This action contains the dynamics of the topological variable in the form of a free rotator
\[ W_Z = \int dt \frac{\dot{N}^2 I}{2}; \quad I = \int \frac{d^3x(D_{i}^{ac}(\Phi_k)\Phi_{0}^{c})^2}{2} = \frac{4\pi}{g^2(2\pi)^2\epsilon} , \] (31)
where $\epsilon$ is a size of the BPS monopole considered as a parameter of the infrared regularization which disappears in the infinite volume limit. The dependence of $\epsilon$ on volume can be chosen so that the density of energy was finite. In this case, the U(1) anomaly can lead to additional mass of the isoscalar meson due to its mixing with the topological variable [18].

The vacuum wave function of the topological free motion in terms of the Pontryagin index (27) takes the form of a plane wave $\exp(iP_N\nu[A^*])$. The well-known instanton wave function [20] appears for nonphysical values of the topological momentum $P_N = \pm i8\pi^2/g^2$ that points out the possible status of instantons as nonphysical solutions with the zero energy in Euclidean space-time $\mathbb{R}^4$. In any case, such the Euclidean solutions cannot describe the phenomena of the type of the complete destructive interference (29).

The Feynman path integral for the obtained unconstrained system in the class of functions of the topological transformations takes the form (see [18])

$$Z_F[l^{(0)}, J^{**}] = \int DN(t) \int \prod_{c=1}^{c=8} [d^2A^c d^2E^{cc*}] \exp \left\{ iW_{l^{(0)}}[A^*, E^*] + i \int d^4x [J^{**}_{\mu} \cdot \hat{A}_\mu^{**}] \right\}, \quad (32)$$

where $J^{cc*}$ are physical sources.

The perturbation theory in the sector of local excitations is based on the Green function $1/D^2(\Phi)$ as the inverse differential operator of the Gauss law which is the non-Abelian generalization of the Coulomb potential. As it has been shown in [18], the non-Abelian Green function in the field of the Wu-Yang monopole is the sum of a Coulomb-type potential and a rising one. This means that the instantaneous quark-quark interaction leads to spontaneous chiral symmetry breaking [8, 21], goldstone mesonic bound states [8], glueballs [21, 22], and the Gribov modification of the asymptotic freedom formula [22]. If we choose a time-axis $l^{(0)}$ along the total momentum of bound states [8] (this choice is compatible with the experience of QED in the description of instantaneous bound states), we get the bilocal generalization of the chiral Lagrangian-type mesonic interactions [8].

The change of variables $A^*$ of the type of (2) with the non-Abelian Dirac factor

$$U(A) = U_Z \exp \left\{ \frac{1}{D^2(\Phi)} D_j(\Phi) \hat{A}_j \right\} \quad (33)$$

and the change of the Dirac dressed sources $J^*$ can remove all monopole physics, including confinement and hadronization, like similar changes (2), (6) in QED (to get a relativistic form of the Feynman path integral) remove all electrostatic phenomena in the relativistic gauges.

The transition to another gauge faces the problem of zero of the FP determinant $detD^2(\Phi)$ (i.e. the Gribov ambiguity [23] of the gauge [26]). It is the zero mode of the second class constraint. The considered example (32) shows that the Gribov ambiguity (being simultaneously the zero mode of the first class constraint) cannot be removed by the change of gauge as the zero mode is the inexorable consequence of internal dynamics, like the Coulomb field in QED. Both the zero mode, in QCD, and the Coulomb field, in QED, have nontrivial physical consequences discussed above, which can be lost by the standard gauge-fixing scheme.

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4 Instead of Conclusion

The variational methods of describing dynamic systems were created for the Newton mechanics. All their peculiarities (including time initial data, spatial boundary conditions $O(1/r)$, time evolution, spatial localization, the classification of constraints, and equations of motion in the Hamiltonian approach) reflect the choice of a definite frame of reference distinguished by the axis of time $t^{(0)} = (1, 0, 0, 0)$. This frame determines also the ”equivalent unconstrained system” for the relativistic gauge theory. This ”equivalent system” is compatible with the simplest variational methods of the Newton mechanics. The manifold of frames corresponds to the manifold of ”equivalent unconstrained systems”. The relativistic invariance means that a complete set of physical states for any ”equivalent system” coincides with the one for another ”equivalent system” [24].

This Schwinger’s treatment of the relativistic invariance is confused with the naive understanding of the relativistic invariance as independence on the time-axis of each physical state. The latter is not obliged, and it can be possible only for the QFT description of local elementary excitations on their mass-shell.

For a bound state, even in QED, the dependence on the time-axis exists. In this case, the time-axis is chosen to lie along the total momentum of the bound state in order to get the relativistic covariant dispersion law and invariant mass spectrum. This means that for the description of the processes with some bound states with different total momenta we are forced to use also some corresponding ”equivalent unconstrained systems”. Thus, a gauge constrained system can be completely covered by a set of ”equivalent unconstrained systems”. This is not the defect of the theory, but the method developed for the Newton mechanics.

What should we choose to prove confinement and compute the hadronic spectrum in QCD: ”equivalent unconstrained systems” obtained by the honest and direct resolving constraints, or relativistic gauges with the lattice calculations in the Euclidean space with the honest summing of all diagrams that lose from the very beginning all constraint effects?

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