CALCULATION OF THE VACUUM ENERGY DENSITY AND GLUON CONDENSATE WITHIN ZERO MODES ENHANCEMENT MODEL OF THE QCD VACUUM

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Abstract

The nonperturbative vacuum structure which emerges from the zero modes enhancement (ZME) model of the true QCD vacuum, appears to be well suited to describe quark confinement, dynamical chiral symmetry breaking (DCSB), current-effective (dynamical)-constituent, as well as constituent-valence quark transformations, the Okubo-Zweig-Iizuka (OZI) rule, dimensional transmutation, etc. It is based on the solution to the Schwinger-Dyson (SD) equation for the quark propagator in the infrared (IR) domain. The importance of the instanton-type fluctuations in the true QCD vacuum for the ZME model is also discussed. This allows to calculate new, more realistic values for the vacuum energy density (apart from the sign, by definition, the bag constant) and the gluon condensate. Our numerical results for the gluon condensate for different numbers of quark flavor \( N_f \) are 2-3 times larger than it is estimated in the QCD sum rules approach. This is in good agreement with recent phenomenological estimates of this quantity.

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I. INTRODUCTION

Today there are no doubts left that the dynamical mechanisms of quark confinement and dynamical chiral symmetry breaking (DCSB) are closely related to the complicated topological structure of the QCD nonperturbative vacuum [1-3]. For this reason, any correct nonperturbative model of quark confinement and DCSB necessary turns out to be a model of the true QCD vacuum and the other way around. Also it becomes clear that the nonperturbative infrared (IR) divergences, are closely related, on one hand, to the above mentioned nontrivial vacuum structure, on the other hand, they are important as far as the large scale behaviour of QCD is concerned [1-5]. If it is true that QCD is an IR unstable theory (has no IR stable fixed point) then the low-frequency modes of the Yang-Mills fields should be enhanced due to the nonperturbative IR divergences. So the gluon propagator can diverge faster than the free one at small momentum, in accordance with $D_{\mu\nu}(q) \sim (q^2)^{-2}$ at small $q$, which describes the zero modes enhancement (ZME) effect in QCD (see our preprint [6] and references therein). If, indeed the low-frequency components of the virtual fields in the true vacuum have a larger amplitude than those of the bare (perturbative) vacuum [4], then the Green’s function for a single quark should be reconstructed on the basis of this effect. One of the main features of this reconstruction is, of course, the correct treatment of this strong singularity within the distribution theory [7] that was precisely done in our previous publications [6, 8]. Let us mentioned in this connection an important observation that the correct treatment of this behaviour within the distribution theory effectively transforms the initial strong singularity ($q^{-4}$) into a Coulomb-like behaviour ($\sim q^{-2}$) in the intermediate and ultraviolet (UV) regions, which is compatible with asymptotic freedom. In fact, this transformation clearly shows that interaction at short distances in QCD within the ZME effect is also completely different from that of quantum electrodynamics (QED). The possible ZME effect was our primary dynamical assumption. We considered this effect as a very similar confining ansatz for the full gluon propagator in order to use it as input information for the quark, ghost Schwinger-Dyson (SD) equations as well as for the corresponding
Slavnov-Taylor (ST) identities [6, 8, 9].

In our recent publication [10] and preprints [6, 11] the basic chiral QCD parameters (the pion decay constant in the current algebra (CA) representation, $F_{CA}$, the quark condensate, the dynamically generated quark mass, $m_d$, etc) have been calculated from first principles within the ZME model of quark confinement and DCSB. As it was mentioned above it was based on the solution to the SD equation for the quark propagator in the infrared (IR) domain [6, 8-11]. At low energies QCD is governed by $SU_L(N_f) \times SU_R(N_f)$ chiral symmetry ($N_f$ is the number of different flavors) and its dynamical breakdown in the vacuum to the corresponding vectorial subgroup [12]. Thus to understand chiral limit physics means to correctly understand the dynamical structure of low energy QCD. A realistic calculation of various physical quantities in this limit becomes important. It is the goal of this work to generalize our calculations to $N_f$ light quarks, as well as to discuss some dynamical aspects of the ZME model which clearly shows the rich possibilities of this effect in QCD.

There are only five independent quantities by means of which all other chiral QCD parameters can be calculated in our model [10], namely

\[
F^2_{CA} = \frac{3}{8\pi^2} k_0^{-2} z_0^{-1} \int_0^{z_0} dz \frac{zB^2(z_0, z)}{\{zg^2(z) + B^2(z_0, z)\}},
\]

\[
m_d = k_0\{z_0B^2(z_0, 0)\}^{-1/2},
\]

\[
\langle \bar{q}q \rangle_0 = -\frac{3}{4\pi^2} k_0^{-3/2} z_0^{-3/2} \int_0^{z_0} dz zB(z_0, z),
\]

\[
\epsilon_q = -\frac{3}{8\pi^2} k_0^{3} z_0^{-2} \int_0^{z_0} dz \{\ln z [zg^2(z) + B^2(z_0, z)] - 2zg(z) + 2\},
\]

\[
\epsilon_g = -\frac{1}{\pi^2} k_0^{-4} z_0^{-2} \times \left[18 \ln(1 + \frac{z_0}{6}) - \frac{1}{2} z_0^2 \ln(1 + \frac{6}{z_0}) - \frac{3}{2} z_0\right],
\]

where $\epsilon_q$ and $\epsilon_g$ describe a single confining quark and nonperturbative gluons (due to the ZME effect) contributions to the vacuum energy density, respectively. The explicit expressions for the solution of the quark SD equation $g(z)$ and $B^2(z_0, z)$ are...
\[ g(z) = z^{-2}[\exp(-z) - 1 + z] \] (1.6)

and

\[ B^2(z_0, z) = 3\exp(-2z) \int_{z_0}^{z} \exp(2z')g^2(z') \, dz' \] (1.7)

respectively. Thus our calculation scheme is self-consistent because we calculate \( n = 5 \) independent physical quantities by means of only \( m = 2 \) free parameters, which possess clear physical sense, so the condition of self-consistency \( n > m \) is well satisfied. The mass scale parameter \( k_0 \) characterizes the region where confinement, DCSB and other nonperturbative effects are dominant while the second independent parameter \( z_0 \) is the constant of integration of the quark SD equation. The details of our scale-setting schemes for system (1.1-1.7) one may find in Refs. [6, 10, 11] and numerical results in the case of a single confining quark from Tables 1 and 2 therein.

II. A POSSIBLE DYNAMICAL PICTURE OF QUARK CONFINEMENT AND DCSB WITHIN THE ZME EFFECT IN QCD

Let us make now a few detailed remarks (some of them are necessary of semi-intuitive as well as of semi-speculative character) shining light on our understanding of the actual dynamical mechanism of quark confinement and DCSB might be interpreted with the help of the possible ZME effect in QCD. Susskind and Kogut [3] have noticed that ”the absence of quarks and other colored objects could only be understood in terms of an IR divergences in the self-energy of a color bearing object”. In our approach gluons remain massless, only zero modes are enhanced. We will discuss in more detail contributions to the self-energy of the colored quark leading first to the dynamical and then to constituent quarks in the context of the ZME effect which is caused by the nonperturbative IR divergences in the true vacuum of QCD.

In order to clarify the dynamical picture which lies at the heart of our model, let us introduce, following Mandelstam [4], two sorts of gluons. The actual (external) gluons are
those which are emitted by a quark and absorbed by another one, while the virtual (internal) gluons are those which are emitted and absorbed by the same quark. Of course, both sorts of gluons are not only free ones. All possible self-energy insertions are assumed to be taken into account as well. At first sight this separation seems to be a simple convention but we will show below that it has a firm dynamical ground, thus it makes our understanding of the above mentioned picture more transparent.

Let us consider now all the possible contributions to the self-energy of a single quark. The most simplest one is shown in Fig. 1. Let us recall that the same self-energy diagram occurs also in the quantum electrodynamics (QED). In contrast to QED, there is an infinite number of additional contributions to the self-energy of a single quark because of the non-abelian nature of QCD, i.e. because of the direct interaction between virtual gluons which is absent in QED. Some of these are shown in Fig. 2. So, from the point of view of the contributions to the self-energy of a single quark, the zero modes are indeed enhanced in QCD in comparison with the electron self-energy in QED. The self-interaction of virtual gluons alone removes a single quark from the mass-shell, making it an effective (dynamical in the chiral limit) object. This is the context of the SD equation describing propagation of a single quark (at large distances) in the true QCD vacuum.

But this is not the whole story yet in QCD because up to now we took into account only contributions induced by the virtual gluons alone. The actual gluons emitted by one quark can contribute to the self-energy of another quark and vice versa. The simplest diagrams of this process are shown in Fig. 3. Moreover contributions shown in Fig. 4 are also inevitable and they describe the process of the conversion (transformation) of virtual gluons into actual ones and the other way around. Thus we consider diagrams, of these type, not as corrections to the cubic and quartic gluon vertices but rather as additional contributions to the self-energy of the quarks. Contributions to the self-energy of each quark will be essentially enhanced in the presence of another quark. In other words each quark additionally enhances the interaction with the vacuum (zero modes) of another quark. Precisely this enhancement of the zero modes by virtue of self-interaction of virtual (internal)
and actual (external) gluons is effectively correctly described by the ZME model.

It is quite plausible that, at large distances between quarks, actual gluons emitted by each quark do repeatedly succeed to convert into virtual ones and vice versa. This leads to a multiple enhancement of the zero modes of each quark. Exactly these additional contributions to the self-energy of each effective quark makes it a constituent object in our model. The mass of the constituent (heavy or light) quark becomes the sum of three terms, namely

\[ m_q = m_{\text{eff}} + \Delta = m_0 + m_d + \Delta. \]  

(2.1)

All terms on the right hand side have clear physical sense. The first term is, obviously, the current mass of a single quark. The second one \( m_d \) describes contributions to the constituent quark mass induced by the self-interaction of virtual gluons alone, while the third term, \( \Delta \), describes contributions to the constituent quark mass which come from the process of the conversion of actual gluons into virtual ones as it was discussed above. This way our model provides a natural dynamical foundation of the current-effective (dynamical)-constituent transformation [13] of the quark degrees of freedom on the basis of non-abelian character of the gluon fields. The existence of a nonzero \( \Delta \) is principal for our model but numerically it should not be large, even for light quarks. Our intuition (based on the obtained numerical results, see Tables 1 and 2 in Refs. [6, 10, 11]) tells us that it is only of the order of a few per cent of the displayed there values of the dynamically generated quark masses, \( m_d \).

In reality the contributions are so mixed up that they can not be separated from each other. There exists an infinite number of possible, topologically complicated, configurations of the vacuum fluctuations of the non-abelian gauge (gluon) fields contributing to the self-energy of each quark while making them constituent objects. The true vacuum of QCD, however, is not settled by these fluctuations alone, its structure is much more richer [14] than that (see also the discussion below). Certainly, a finite number of favorable, topologically distinct vacuum configurations, which minimize the energy of the bound states, should exist. It is hard to believe that in the real word of four dimensions, the favorable topologically complicated configurations are strings or planar ones. In this context, it is im-
portant to comprehend that the linearly rising quark-antiquark potential at large distances, nicely showed by recent lattice calculations [15], is not a privilege of the planar or string configurations only. Though the above mentioned linear potential does not contradict the ZME effect, nevertheless the potential concept in general is a great simplification of the real dynamical picture which emerges from our model. As it was underlined in our papers [6, 8-11], the enhancement of zero modes necessary leads to full vertices while the potential concept of the constituent quark model (CQM) is based on point-like ones. Moreover, we think that the potential concept has already played its useful role and now should be retired from the scene like the Bohr orbits after the creation of the true theory of atoms - quantum mechanics.

It is plausible that these energetically advantageous configurations of vacuum fluctuations (leading to the formation of the bound-states of the constituent quarks) occur at a certain distance between constituent quarks. They will be completely deformed (or even destroyed) if one attempts to separate the constituent quarks further from each other. The nonperturbative vacuum of QCD is filled with quark-antiquark virtual pairs which consist of various components of quark degrees of freedom (light, heavy, constituent, dynamical, etc). This is an inevitable consequence of the ZME effect. As a result of the above mentioned nontrivial topological deformation, at least one quark loop will be certainly "cut". We may, for convenience, think of this as such a topological deformation which allows for quarks from the loop to recombine with the initial constituent quarks. It is evident that the breaking of the gluon line is not so important as the above mentioned cut of the quark loop which is equivalent to the creation of the corresponding quark-antiquark pairs from the vacuum. The vacuum of QCD will be immediately rearranged and, instead of "free" constituent quarks, new hadron states will occur.

At short distances the situation is completely different from the above described. Indeed, in this case the actual gluons emitted by each quark do not repeatedly succeed to convert into virtual ones. So, at these distances, interaction between quarks proceeds mainly through the exchange of actual gluons. This means that the interaction of the constituent quarks with
the vacuum (i.e. contributions to its self-energy), due to the above mentioned process of conversion, is essentially decreased. So they become valence quarks. The intensity of the process of conversion determines the constituent-valence transformation. One may say that the constituent (valence) quarks are ”valence (constituent)-in-being quarks”. In other words, if the process of conversion becomes stronger then the valence quark becomes constituent and vice versa. If the process of conversion becomes weaker then the constituent quark becomes valence, so hadron becomes consisting of valence and sea quarks and mainly actual gluons. Precisely this picture of hadrons emerges from deep inelastic scattering experiments.

A. The Okubo-Zweig-Iizuka selection rule

Any correct model of quark confinement should explain at least qualitatively the famous Okubo-Zweig-Iizuka (OZI) selection rule [16] since it reflects the unusual dynamics of quarks inside hadrons, that, in turn, is closely related to the QCD vacuum structure. In the context of the ZME model it becomes clear that the topological rearrangement of the vacuum by means of the direct annihilation of the initial (final) constituent quarks, entering the same hadrons, is hardly believable. In fact, what does the above mentioned direct annihilation mean? This would mean that the initial (final) constituent quarks, emitting (absorbing) a number of gluons, can annihilate with each other without the break-up of the corresponding quark loops in our model. The initial (final) constituent quarks always emit and absorb gluons in each preferable configuration. Nothing interesting should happen during these processes. This is a normal phase of each preferable configuration and it describes only its trivial rearrangement. Any nontrivial rearrangement of the vacuum can only begin with cutting the quark loop. As it was mentioned above, this is equivalent to the creation of a quark-antiquark pair from the vacuum. Then the annihilation of the initial (final) constituent quarks with the corresponding quarks, liberated from the loop, becomes possible. Diagramatically this looks like a direct annihilation (see Fig. 5). The probability to create the necessary pair, in order to annihilate the initial (final) constituent quarks, is rather small,
so, in general, the annihilation channel must be suppressed. It is much more probable for the quarks liberated from the loop to recombine with the initial (final) constituent quarks in order to generate new hadron states. In more complicated cases (when many quark loops are cut) the annihilation of the initial (final) constituent quarks becomes more probable and this process will compete with the process of the recombination of the initial (final) constituent and liberated-from-the-vacuum quarks to generate new hadron states.

The heavy constituent quarks inside hadrons (for example, in $c\bar{c}$ systems) are much closer to each other (the distance between them is of the order $m_h^{-1} \ll m_q^{-1}$, where $m_h$ and $m_q$ denote the masses of the heavy and light constituent quarks, respectively) than their light counterparts. However at short distances the interaction between them proceeds mainly not through the vacuum fluctuations but via the exchange of the actual gluons as this was explained above. This means, in turn, that the number of virtual loops which should be cut is small. Also the probability to cut heavy quark loop, or equivalently to create a heavy quark-antiquark pair from the vacuum, is much less than to create, for example, a light pair. The fluctuations in the density of instantons (see next section) and condensates during the vacuum’s rearrangement also come into play. Thus the process of the rearrangement of the vacuum, leading to the transition between hadrons (their strong decays) on the basis of the annihilation of the initial (final) constituent quarks, as shown in Fig. 5, should be suppressed in comparison with the process of the recombination of the initial (final) constituent quarks with the quarks liberated from the vacuum. This is shown in Fig. 6. Thus in the case of heavy quarks our qualitative dynamical explanation of the OZI rule is in agreement with its standard explanation which argues that the QCD coupling constant becomes weak at short distances and suppresses the annihilation channel. However this argument fails to explain why the violation of the OZI rule for the pseudoscalar octet is bigger than for the vector one.

Let us analyse this problem in our approach. Light constituent quarks inside pseudoscalar and vector mesons are at relatively large distances ($\sim m_q^{-1}$ from each other) than their heavy counterparts, for example in $c\bar{c}$ systems. This means that interaction between them
is mediated mainly by the vacuum fluctuations which provide plenty of various quark loops to be cut during the process of the vacuum rearrangement. So the annihilation channel should not be suppressed for these octets. Indeed, the violation of the OZI rule in the pseudoscalar mesons is not small, but for the vector mesons it is again small, i.e. comparable to the violation in the $c\bar{c}$ systems. Our model provides the following explanation for this problem: The same quark-antiquark pair in pseudoscalar and vector mesons is in the same $S$-state. The only difference between them is in the relative orientation of the quark spins. Quark and antiquark spins are oriented in the same direction in the vector mesons, while in the pseudoscalar mesons their orientation is opposite. This is schematically shown in Figs. 7 and 8. For light constituent quarks spin effects become important, while for heavy constituent ones such a relativistic effect as spin and, in particular, its orientation is not so important. As it was repeatedly mentioned above, a nontrivial rearrangement of the vacuum in our model always starts from the cut of the quark loops. In order to analyse the violation of the OZI rule, from the point of view of annihilation of spin degrees of freedom, let us think of quark loops as "spin loops". In pseudoscalar mesons at least one spin-antispin liberated-from-the-vacuum-pair is needed to annihilate the initial pair. This is schematically shown in Fig. 7. In vector mesons at least two spin-antispin pairs are needed for this purpose, but, in addition, an intermediate meson (exited) state certainly appears, see Fig. 8. This means that the annihilation channel for the vector mesons, unlike for the pseudoscalars ones, is suppressed. It is worth noting that the OZI rule is a selection rule and it is not a conservation law of some quantum number. So its breakdown is always possible and suppressed processes may proceeds through the appropriate intermediate states [17]. Exactly this is shown in Fig. 8 schematically. All this explains the violation of the OZI rule in the pseudoscalar channel in comparison with the vector one in our model. For vector mesons the decay $\phi \to 3\pi, \rho\pi$, proceeding though the annihilation channels, are suppressed in comparison with the decay $\phi \to K^+K^-$ which occurs via the recombination channel while for pseudoscalar mesons, say, the decay $\phi \to 3\pi$ is not suppressed. In order to confirm our qualitative explanation of the OZI selection rule quantitatively, it is necessary to calculate
the strong decay widths of mesons in our model, which, of course, is beyond the scope of
the present paper.

III. INSTANTONS

The main ingredients of the QCD vacuum, in our model, are quark and gluon condensates, quark-antiquark virtual pairs (sea quarks) and self-interacting nonperturbative gluons. The vacuum of QCD has, of course, much more remarkable (richer) topological structure than this. It is a very complicated medium and its topological complexity means that its structure can be organized at various levels and it can contain perhaps many other components [1, 14] besides the above mentioned. There are a few models of the nonperturbative vacuum of QCD which are suggestive of what a possible confinement mechanism might be like (see recent paper [18] and references therein). We will not discuss these models; let us only mentioned that the QCD-monopole condensation model proposed by t’ Hooft and Mandelstam [19] within the dual Ginzburg-Landau effective theory [20] also invokes the ZME effect as well as the mechanism (classical) of the confining medium recently suggested by Narnhofer and Thirring [21]. Let us ask one of the main questions now.

What is a mechanism like which initiates a topologically nontrivial rearrangement of the vacuum? It is already known, within our model, that this may begin with the cut, at least, of one quark loop and therefore, at least, one quark-antiquark pair emerges from the vacuum. Why can the quark loop be cut at all and what prevents the quarks from the cut loop to annihilate again with each other? The fluctuations in the nonperturbative vacuum of QCD must exist which do these job. This is an inevitable consequence of our model of the vacuum. We see only one candidate for this role in four dimensional QCD, namely instantons and anti-instantons and their interactions [1, 14, 22].

Instantons are classical (Euclidean) solutions to the dynamical equation of motion of the nonabelian gluon fields and represent topologically nontrivial fluctuations of these fields. Self-interaction of gluons should be important for the existence of the instanton-like fluc-
tuations even at classical level. In the random instanton liquid model (RILM) [23], light quarks can propagate over large distances in the QCD vacuum by simply jumping from one instanton to the next one. In contrast, in our model the propagation of all quarks is determined by the corresponding SD equations (due to the ZME effect) so that they always remain off mass-shell. Thus we do not need the picture of jumping quarks. As opposed to the RILM, we think that the main role of the instanton-like fluctuations is precisely to prevent the quarks and gluons from freely propagating in the QCD vacuum. Running against instanton-like fluctuations, the quarks undergo difficulties in their propagation in the QCD vacuum which, as was mentioned above, is a very complicated inhomogeneous medium. At some critical value of the instanton density the free propagation of the virtual quarks from the loops become impossible so they never annihilate again with each other. Obviously, this is equivalent to the creation of the quark-antiquark pairs from the vacuum. From this moment the nontrivial rearrangement of the vacuum may start. The role of the instanton-like fluctuations appears to be "cutting" the quarks loops and preventing them from the immediate annihilation of quarks and antiquarks liberated from the loops. Thus liberated from loops quarks may, in principle, recombine or even annihilate with initial constituent quarks to produce new hadron states. In this way precisely instantons may promote transitions between hadrons, i.e. they destabilize energetically advantageous (dominant) configurations of the vacuum fluctuations which lead to hadron states. One of the main features of the instanton-induced effects is tunneling between topologically distinct vacuums in Minkovski space [1]. Our understanding of their role in the QCD vacuum structure is in agreement with this.

Being classical (not quantum!) fluctuations, instantons can cut the quark loops in any points even in the quark-gluon vertices. A simple cut of the quark loops, as it was repeatedly emphasized above, is equivalent to the creation of the corresponding quark-antiquark pairs from the vacuum. Let us ask now what happens if all the external (actual) gluon lines will be cut from the quark loops by the instantons. Well, in this case each quark loop becomes a closed system. Because of the vacuum pressure they immediately collapse and
one obtains nothing else but quark condensate if, of course, all internal gluon lines can be included into the quark self-energy. Another scenario is also possible when not all the internal gluon lines can be included into the quark or gluon self-energy. In general, the presence of the internal gluon lines in the vacuum diagrams prevent them from collapsing (because they counterbalance the vacuum pressure) and consequently they should contribute to the vacuum energy density.

It would be suggestive to conclude that the same mechanism works in order to produce gluon condensates despite the much more complicated character of the gluon self-interactions. But this is not the case indeed, since there is a principal difference between quark and gluon condensates. The former ones do not contribute to the vacuum energy density (in the chiral limit, see below) and, in this sense, they play the role of some external field. While the latter ones, as was shown first in Ref. [24], are closely related to the vacuum energy density. Like we noted previously, nothing interesting should be happen if the instantons cut the gluon line in one point. Let us imagine that many gluon lines will be cut by the instantons in many points. The vacuum becomes filled up with gluon pieces (segments) which, due to the existence of the gluon strong self coupling, can recombine, in principle, as some colour singlet bound-states – gluonia or glueballs. Unlike quark condensates, glueballs should have internal pressure because of the strong self-interaction of composite gluon segments which prevents them from collapsing. In turn, this means that the glueballs should be heavy enough.

The pseudoscalar mesons (consisting of light quarks) are Nambu-Goldstone (NG) states so their masses remain zero in the chiral limit even in the presence of DCSB. From our model it follows that the existence of the instanton-type fluctuations in the true vacuum promote strong meson decays by preventing quarks and gluons from freely propagating in it. So one can conclude in that instanton-type fluctuations should be totally suppressed in this case in order to provide stability for the massless NG states since massless particles cannot decay. This feature of the instanton physics in the massless quarks case was discovered by t’ Hooft [25]. In the presence of DCSB, however, the instanton-type fluctuation are restored [14,
26], but the contribution of the instanton component to the vacuum energy density, $\epsilon_I$, still remains small. So the dilute gas approximation for the instanton component seems to be relevant in this case. Not going into details of the instanton physics (well described by Callan, Dashen and Gross in Ref. [14]), let us only emphasize that at short distances the density of small size instantons should rapidly decrease and conversely increase at large distances where large scale instantons, anti-instantons and their interactions also come into play. Otherwise it would be difficult to understand the role which we would like to assign to the instantons in our model. This is in agreement with the behaviour of the instanton component of the QCD vacuum at short and large distances described in the above mentioned paper [14], as well as with our understanding of the actual dynamical mechanism of quark confinement and DCSB. It is possible to say that we treat the instanton component of the QCD vacuum not as a "liquid" but rather as a "forest" in which instantons and anti-instantons are considered as "trees" with $SU_c(3)$ orientation, position and scale size.

IV. CALCULATION OF THE VACUUM ENERGY DENSITY

The nontrivial rearrangement of the vacuum can start only when the density of instantons achieves some critical values, different for all distinct vacuums. For this reason, despite being a classical phenomena, instantons should nevertheless contribute to the vacuum energy density through the above mentioned quantum tunneling effect which is known to lower the energy of the ground-state. As it was noticed above, $\epsilon_I$ should be small for light quarks with dynamically generated masses. However, the same conclusion seems to be valid for heavy quarks as well in our model. Heavy quarks are at short distances from each other (in mesons) at which the nonperturbative effects, such as instantons and enhancement of zero modes, are suppressed. Thus the dilute gas approximation for the instanton component seems to be applicable to light quarks with dynamically generated masses as well as to heavy quarks. In this case it is worth assuming, following the authors of Ref. [24], that light and heavy quarks match smoothly (in our model this is almost inevitable consequence). In the
above mentioned RILM [23] of the QCD vacuum, for a dilute ensemble, one has
\begin{align}
\epsilon_I &= -(1/4)(11 - \frac{2}{3} N_f) \times 1.0 \text{ fm}^{-4} \\
&= -(0.00417 - 0.00025 N_f) \text{ GeV}^4. \tag{4.1}
\end{align}

For $N_f = 3$ it coincides with the estimate of the QCD sum rules approach on account of the phenomenological value of the gluon condensate [24] (via the trace anomaly relation, see below). It decreases with increasing $N_f$ that defies a physical interpretation of the vacuum energy density as the energy per unit volume. That is why the value of the vacuum energy density at the expense of the instanton contributions alone is at least not complete.

In QCD the vacuum energy density (at the fundamental quarks and gluons level) should be calculated through the effective potential method for composite operators [27] (for review see Refs. [28]), since in the absence of the external sources the effective potential is nothing but the vacuum energy density. This method gives it in terms of loop expansion in powers of Plank constant. Using it we have already calculated [6, 10] the contributions to the vacuum energy density of confining quarks with dynamically generated masses, $\epsilon_q$ (1.4), and nonperturbative gluons, $\epsilon_g$ (1.5), due to enhancement of zero modes (both at log-loop level). Thus the value of the vacuum energy density, as given by the ZME model, is
\begin{equation}
\epsilon_{ZME} = \epsilon_g + N_f \epsilon_q, \tag{4.2}
\end{equation}

where we introduced the dependence on the number of different quark flavors $N_f$ since $\epsilon_q$ itself is the contribution of a single confining quark. Numerically it is
\begin{equation}
\epsilon_{ZME} = -(0.0016 + 0.0015N_f) \text{ GeV}^4. \tag{4.3}
\end{equation}

(For numerical values of each component $\epsilon_g$ and $\epsilon_q$ in calculation scheme A see our paper [10] (Table 1)).

However, neither contribution (4.1) nor (4.3) is complete. It was already explained in detail above why the instanton-type fluctuations are needed for the ZME model. In order to get a more realistic value of the vacuum energy density, let us add $\epsilon_I$, as given by Eq. (4.1),
to the ZME model value (4.3), i.e. let us put that the total (t) vacuum energy density at least is

$$\epsilon_t = \epsilon_I + \epsilon_{ZME} = \epsilon_I + \epsilon_g + N_f \epsilon_q,$$

(4.4)

and numerically it is

$$\epsilon_t = -(0.00577 + 0.00125N_f) \text{GeV}^4.$$

(4.5)

Note, the dependence on $N_f$ now becomes correct. The above described components produce the main (leading) contribution to the vacuum energy density. The next-to-leading contributions (as given by the effective potential for composite operators at two-loop level [27]) are $h^2$-order, where $h$ is the above mentioned Plank constant. Thus they are suppressed at least by one order of magnitude in comparison with our values (4.5).

In the calculation scheme B bounds for the total vacuum energy are

$$-0.00638 - 0.00111N_f \leq \epsilon_t \leq -0.00461 - 0.00145N_f$$

(4.6)

in units of $\text{GeV}^4$. (For numerical values of each component $\epsilon_g$ and $\epsilon_q$ in calculation scheme B see our paper [10] (Table 2)).

V. CALCULATION OF THE GLUON CONDENSATE

The vacuum energy density is important in its own right as the main characteristics of the nonperturbative vacuum of QCD. Furthermore it assists in estimating such an important phenomenological parameter as the gluon condensate, introduced within the QCD sum rules approach to resonance physics [24]. Indeed, because of the Lorentz invariance,

$$\langle 0 | \Theta_{\mu\nu} | 0 \rangle = 4\epsilon_t$$

(5.1)

holds where $\Theta_{\mu\nu}$ is the trace of the energy-momentum tensor and $\epsilon_t$ is the sum of all possible contributions to the vacuum energy density. According to QCD the famous trace anomaly relation [29] in the general case (nonzero "bare" quark masses $m_f$) is
\[ \Theta_{\mu\nu} = \frac{\beta(\alpha_s)}{4\alpha_s} G_{\mu\nu}^a G_{\mu\nu}^a + \sum_f m_f \bar{q}_f q_f. \] (5.2)

(G_{\mu\nu}^a being the gluon field strength tensor). The function \( \beta(\alpha_s) \), up to terms of order \( \alpha_s^3 \), is [1]

\[ \beta(\alpha_s) = -(11 - \frac{2}{3} N_f) \frac{\alpha_s^2}{2\pi} + O(\alpha_s^3). \] (5.3)

Sandwiching Eq. (5.2) between vacuum states and on account of relations (5.1) and (5.3), one obtains

\[ \epsilon_t = -(11 - \frac{2}{3} N_f) \frac{1}{32} \langle 0 | \frac{\alpha_s}{\pi} G_{\mu\nu}^a G_{\mu\nu}^a | 0 \rangle + \frac{1}{4} \sum_f m_f \langle 0 | \bar{q}_f q_f | 0 \rangle, \] (5.4)

where \( \langle 0 | \bar{q}_f q_f | 0 \rangle \) is the quark condensate and \( \langle 0 | \frac{\alpha_s}{\pi} G_{\mu\nu}^a G_{\mu\nu}^a | 0 \rangle \) is nothing but the gluon condensate [24]. The weakness of this derivation is, of course, relation (5.3) which holds only in the perturbation theory. In any case, it would be enlightening to numerically estimate the gluon condensate with the help of relation (5.4).

From Eqs. (4.5) and (5.4) in the chiral limit \( (m_f = 0) \) it finally follows

\[ \langle 0 | \frac{\alpha_s}{\pi} G_{\mu\nu}^a G_{\mu\nu}^a | 0 \rangle = \frac{0.00577 + 0.00125 N_f}{0.344 - 0.021 N_f} \text{ GeV}^4. \] (5.5)

Numerically our values are:

\[ \langle 0 | \frac{\alpha_s}{\pi} G_{\mu\nu}^a G_{\mu\nu}^a | 0 \rangle (N_f = 0) \simeq 0.01677 \text{ GeV}^4, \]
\[ \langle 0 | \frac{\alpha_s}{\pi} G_{\mu\nu}^a G_{\mu\nu}^a | 0 \rangle (N_f = 1) \simeq 0.0217 \text{ GeV}^4, \]
\[ \langle 0 | \frac{\alpha_s}{\pi} G_{\mu\nu}^a G_{\mu\nu}^a | 0 \rangle (N_f = 2) \simeq 0.0274 \text{ GeV}^4, \]
\[ \langle 0 | \frac{\alpha_s}{\pi} G_{\mu\nu}^a G_{\mu\nu}^a | 0 \rangle (N_f = 3) \simeq 0.034 \text{ GeV}^4. \] (5.6)

Thus our values of the gluon condensate (5.5-5.6) are of factor of 2-3 larger than the value phenomenologically estimated from the QCD sum rules approach [24], namely

\[ \langle 0 | \frac{\alpha_s}{\pi} G_{\mu\nu}^a G_{\mu\nu}^a | 0 \rangle \simeq 0.012 \text{ GeV}^4. \] (5.7)

There exist already phenomenological estimates of the gluon condensate pointing out that its so-called standard value, which is determined by the instanton-type fluctuations alone
(5.7), is too small. It was pointed out (perhaps first) in Ref. [30] (see also Ref. [31]) that QCD sum rules substantially underestimated the value of the gluon condensate about of factor of 2-3. Our numerical results (5.5-5.6) are in good agreement with these estimate. The most recent phenomelogical calculation of the gluon condensate is given by Narison in Ref. [32], where a brief review of many previous calculations is also given. His analysis leads to the update average value as

\[
\langle 0 | \frac{\alpha_s}{\pi} \pi G^a_{\mu\nu} G^a_{\mu\nu} | 0 \rangle = (0.0226 \pm 0.0029) \text{ GeV}^4. \tag{5.8}
\]

Thus our results (5.5-5.6), calculated from first principles on the basis of the ZME model, are in good agreement with phenomenologically estimated values of the gluon condensate summarized in the Narison’s paper (see Table 2 in Ref. [32]).

In the calculation scheme B bounds for the gluon condensate on account of Eqs. (4.6) and (5.4) become

\[
\frac{0.00461 + 0.00145 N_f}{0.344 - 0.021 N_f} \leq \langle 0 | \frac{\alpha_s}{\pi} G^a_{\mu\nu} G^a_{\mu\nu} | 0 \rangle \leq \frac{0.00638 + 0.00111 N_f}{0.344 - 0.021 N_f} \text{ GeV}^4 \tag{5.9}
\]

in units of GeV$^4$. Numerically these bounds are:

\[
0.0134 \leq \langle 0 | \frac{\alpha_s}{\pi} G^a_{\mu\nu} G^a_{\mu\nu} | 0 \rangle (N_f = 0) \leq 0.0185,
\]
\[
0.0187 \leq \langle 0 | \frac{\alpha_s}{\pi} G^a_{\mu\nu} G^a_{\mu\nu} | 0 \rangle (N_f = 1) \leq 0.0232,
\]
\[
0.0248 \leq \langle 0 | \frac{\alpha_s}{\pi} G^a_{\mu\nu} G^a_{\mu\nu} | 0 \rangle (N_f = 2) \leq 0.0285,
\]
\[
0.032 \leq \langle 0 | \frac{\alpha_s}{\pi} G^a_{\mu\nu} G^a_{\mu\nu} | 0 \rangle (N_f = 3) \leq 0.0345. \tag{5.10}
\]

These bounds are also in agreement with values of the gluon condensate summarized in the above mentioned Narison’s paper (Table 2 in Ref. [32]). The lower and upper bounds for the gluon condensate

\[
0.04 \leq \langle 0 | \frac{\alpha_s}{\pi} G^a_{\mu\nu} G^a_{\mu\nu} | 0 \rangle \leq 0.105 \text{ GeV}^4, \tag{5.11}
\]

recently derived from the families of $J/\Psi$ and $\Upsilon$ mesons in Ref. [33] slightly and substantially overestimate our lower and upper bounds (5.10), respectively. The tendency, however, to increase the value of the gluon condensate is correct.
VI. THE BAG CONSTANT

A nontrivial relation between our model, on one hand, and the bag [34] and string [35] models, on the other hand, would not be surprised. In this connection some dynamical aspects of our model should be underlined. From the above consideration it follows that, from a dynamical point of view, maybe the ZME effect does not lead to string configurations of flux tube type between quarks. Nevertheless, there is no doubt that this dynamical process works like a string preventing quarks to escape from each other. It takes place in the finite volume of the QCD vacuum but it does not require the introduction of an explicit surface. The finiteness of the cut-off $z_0$ results in unphysical singularities (at this point $z_0$) of the solutions to the quark SD equation which are due to inevitable ghost degrees of freedom in QCD. This has nothing to do with the bag fixed boundary. Numerically it depends on a scale at which nonperturbative effects become essential in our model. We treat a hadron as a dynamical process which takes place in some finite volume of the vacuum rather than as an extended object with an explicitly fixed surface in the vacuum. The ZME model remains a local field theory. However, the existence of the vacuum energy per unit volume – the bag constant $B$ – is important in our model as well. The inward positive pressure $B$ counterbalances the vacuum energy density needed for generating the vacuum fluctuations, inspired by the enhancement of the zero modes in our model, i.e. the sum of the bag constant and the nonperturbative vacuum energy density must be zero. We consider the bag constant as a universal one which characterises the complex nonperturbative structure of the QCD vacuum itself and it does not depend on the hadron matter.

The bag constant is defined as the difference between the energy density of the perturbative and the nonperturbative QCD vacuums. We normalized the perturbative vacuum to zero [6, 10], so in our notations the bag constant becomes

$$B = -\epsilon_t = (0.00577 + 0.00125N_f) \text{GeV}^4,$$

on account of Eq. (4.5). So our predictions for a more realistic values of the bag constant are:
\( B(N_f = 0) \simeq 0.006 \, GeV^4 \simeq (278 \, MeV)^4 \simeq 0.78 \, GeV/fm^3, \)
\( B(N_f = 1) \simeq 0.007 \, GeV^4 \simeq (290 \, MeV)^4 \simeq 0.91 \, GeV/fm^3, \)
\( B(N_f = 2) \simeq 0.008 \, GeV^4 \simeq (300 \, MeV)^4 \simeq 1.04 \, GeV/fm^3, \)
\( B(N_f = 3) \simeq 0.0095 \, GeV^4 \simeq (312 \, MeV)^4 \simeq 1.24 \, GeV/fm^3. \)  \hspace{1cm} (6.2)

It has been noticed in [36] that nobody knows yet how big the bag constant might be, but generally it is thought it is about 1 GeV/fm\(^3\). The predicted value for the most relevant physical case of \( N_f = 2 \) is in fair agreement with this expectation.

In the calculation scheme B bounds for the bag constant on account of Eq. (4.6) are
\[
0.00461 + 0.00145 N_f \leq B \leq 0.00638 + 0.00111 N_f
\]  \hspace{1cm} (6.3)
in units of GeV\(^4\).

**VII. DISCUSSION**

Let us make now a few things perfectly clear. It makes sense to underline once more that the vacuum energy density is not determined by the trace anomaly relation (5.4). It is much more fundamental quantity than the gluon condensate and the gluon condensate itself is determined by the vacuum energy density via the trace anomaly relation (5.4) in the chiral limit. As was explained above, the vacuum energy density in QCD is to be calculated completely independently from the gluon condensate. The gluon or quark condensates may or may not exist but the vacuum energy density as energy per unit volume always exist. In QCD at quantum level it should be calculated (as it was pointed out above) by the effective potential method for composite operators [27]. It gives the vacuum energy density in the form of loop expansion where the number of the two-particle irreducible vacuum loops (consisting mainly of confining quark and nonperturbative gluon degrees of freedom) is equal to the power of the Plank constant. So the vacuum energy density at quantum level to leading order becomes in general \( \epsilon_g + N_f \epsilon_q + O(h^2) \). We beleive that the ZME model correctly reproduces these contributions to the vacuum energy density (see Eqs. (4.2-4.3)).
But this is not the whole story in QCD. The instanton-type fluctuations also exist in the nonperturbative vacuum of QCD and they contribute to the vacuum energy density as well. This contribution, however, is a contribution at classical level. So the vacuum energy density becomes the sum of all possible contributions and only this sum determines the gluon condensate in the chiral limit via the trace anomaly relation. That is why the realistic value of the gluon condensate is approximately 2-3 times larger than it is estimated in QCD sum rules.

As it was already mentioned above, a nontrivial rearrangement of the QCD true vacuum will occur if the density of the instanton-like fluctuations achieves some critical value. This critical value can be reached when $-\epsilon_I \gtrsim -(\epsilon_g + \epsilon_q)$, i.e. when at least one sort of quark flavors is presented in the QCD vacuum. Using our numerical results (4.3) as well as (4.1) for $N_f = 1$, it is easy to see that this condition is nicely satisfied, namely $0.00395\, GeV^4 \gtrsim 0.0031\, GeV^4$. Within the ZME model, this clearly shows that Shuryak [23] apparently correctly suggested the average distance between instantons as 1 $fm$ and consequently the density of instantons as $1\, fm^{-4}$ although he started from the underestimated value of the gluon condensate which was phenomenologically given by the QCD sum rules approach (5.7). Then he used the trace anomaly relation (5.4) in the chiral limit for $N_f = 3$ and thus reproduced the value of the vacuum energy density (4.1) which is due to the instanton contributions alone. At that moment nobody knew how to calculate the vacuum energy density at quantum level and the trace anomaly relation was a unique way to estimate it.

Quite recently in quenched ($N_f = 0$) lattice QCD by using the so-called "cooling" method the role of the instanton-type fluctuations in the QCD vacuum was investigated [37]. In particular, it was concluded that the instanton density should be $n = (1 + \delta) \times fm^{-4}$, where $\delta$ was estimated as $\delta \simeq 0.3 - 0.6$ depending on cooling steps. So this enhancement in the density of instantons leads to the enhancement of the vacuum energy density due to the instantons contributions alone, leaving contributions from other components unchanged, of course. Here two scenario are possible. First, this really takes place and Eq. (4.1) should be improved on account of the above mentioned enhancement in the density of instantons. Let
us argue, however, that Eq. (4.1) still remains valid and the enhancement of the vacuum energy density is due to the nonperturbative gluons contributions which apparently also can not be removed by cooling method from the QCD vacuum. Indeed, the additional contribution numerically is

\[ \epsilon_\delta = -(11/4)(0.3 - 0.6) \times f m^{-4} \]
\[ = -(0.00125 - 0.0025) \text{GeV}^4. \]  

(7.1)

The average between these values \( \epsilon_\delta \approx 0.0019 \text{GeV}^4 \) is rather close to our value for \( \epsilon_g \approx 0.0016 \text{GeV}^4 \) (see Eq. (4.3)). If this is so then there is no need to change key parameters of RILM [23] and perhaps of the interacting instanton liquid model (IILM) [38]. What is necessary indeed is to take into account other possible contributions to the vacuum energy density.

Thus one may conclude in that Eq. (4.1) rather correctly gives the contribution of the instanton component to the total vacuum energy density while the value of the gluon condensate is precisely determined by the total vacuum energy density. This is the reason (as was mentioned above) why the realistic value of the gluon condensate 2-3 times larger than the instanton component alone can provide. We have shown that additional contributions to the vacuum energy density and consequently to the gluon condensate are due to the nonperturbative gluons and confining quarks on the basis of the ZME model of the QCD true vacuum. Our numbers for the gluon condensate’s values (obtained from first principles) are in good agreement with recent phenomenological estimates summarized by Narison in Ref. [32].

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FIGURES

FIG. 1. The simplest contribution to the quark self-energy induced by a virtual (internal) gluon.

FIG. 2. The simplest contributions to the quark self-energy induced by the self-interactions of virtual (internal) gluons.

FIG. 3. The simplest contributions to the quark self-energy induced by an actual (external) gluons emitted by another quark.

FIG. 4. The simplest contributions to the quark self-energy due to the processes of the conversion (transformation) of the virtual gluons into the actual ones and vice versa.

FIG. 5. The quark diagram for the decay of the $J/\Psi$ meson. The dashed line schematically shows the $c\bar{c}$ pair emerged from the nonperturbative QCD vacuum. The true number of intermediate gluons is dictated by colour and charge conservation.

FIG. 6. The quark diagram for the decay of the $\Psi''$ meson. The dashed line schematically shows the $u\bar{u}$ pair emerged from the nonperturbative QCD vacuum.

FIG. 7. The quark diagram for the decay of the pseudoscalar (P) particle. The arrows show the reciprocal orientation of the spins. The dashed line schematically shows the pair emerged from the QCD vacuum. The true number of intermediate gluons is dictated by colour and charge conservation.

FIG. 8. The quark diagram for the decay of the vector (V) meson. The arrows show the reciprocal orientation of the spins. The dashed lines schematically show the pairs emerged from the QCD vacuum. The true number of intermediate gluons is dictated by colour and charge conservation. An intermediate meson (exited) state is inevitable.
