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To apply QCD to nuclear physics one needs methods of long-distance QCD. A new method. method of Confining Background Fields, CBF, which incorporates confinement, is presented with applications to heavy and light quarks, both in mesons and baryons. Spin-dependent forces are calculated for light and heavy quarks. The quark potential model in some limiting case is derived.

Fig. - , ref. - 25

1. INTRODUCTION

There is still a wide gap between QCD and models used in hadronic and nuclear physics, such as bag models [1], string models [2], quark potential models, QPM (see [3] and refs.therein). The QPM and bag models have been used for hadron-hadron interaction,and in particular the Quark Compound Bag model (QCB) has been very successful in describing NN and other hadronic systems [4]. However the QCB model contains phenomenological parameters which are not directly deducible from QCD.

The QPM is very successful for single hadrons, e.g. for heavy quarkonia [5], mesons and baryons [3] but it was never derived from QCD and till recently one could not explain why it works in so wide domain of applications.

For multiquark states and, more generally, for hadron-hadron interactions the QPM was not that successful [6] and suffers from some drawbacks, e.g. the Van-der-Waals forces. At the same time an alternative - the flux-tube model has been developed [7], starting from the lattice strong coupling expansion.

All this gives hope that the gap between QCD and hadronic and nuclear physics will be bridged from two directions: the quark hadronic models will be more and more fundamental, and methods of nonperturbative QCD more and more powerful to include finally e.g. the NN interaction. The present talk concerns with the last direction.

2. PROPERTIES OF QCD VACUUM

The QCD vacuum is filled with dense color-field distributions, which are characterized by the gluonic condensate [8]:

$$\frac{\alpha_s}{\pi} \langle F^a_{\mu\nu} F^a_{\mu\nu} \rangle = 0.012 \text{ GeV}^4 \quad (1)$$

so that the average fields are strong $\langle |F| \rangle \sim 0.5 \text{ GeV}^2$ as compared to $\Lambda_{QCD} \sim 0.1 \div 0.15$ GeV. To reconcile nonzero vacuum fields with Lorentz invariance one necessarily arrives at the picture of stochastic distribution of vacuum fields, with color, strength and Lorentz direction randomly changing from point to point. This distribution can be characterized by averaged correlators, e.g.

$$\langle E^a_i(x)\, E^a_k(y) \sim \delta_{ik} D(x-y) + \partial_i [(x-y)_k D_1(x-y)] \quad (2)$$

where we have omitted a phase factor for simplicity and explicitly expressed the correlator in terms of two scalar functions D, D_1. Using the Lorentz (Poincaré) invariance of the vacuum stochastic ensemble, D and D_1 should depend only on $(x-y)^2$.

Higher order correlators are possible, e.g. $\langle F(x_1)\ F(x_2) \ldots F(x_n) \rangle$, and will be considered later. We keep for simplicity only D and D_1 and as will be shown in the next sections, all main features of long-distance dynamics, including confinement, can be obtained already in this approximation. To this end we introduce the CBF method.

3. METHOD OF CONFINING BACKGROUND FIELDS (CBF)

Suppose we know the distribution of all gluonic degrees of freedom in the vacuum. In Monte Carlo simulations this is stored as values of gluon fields at each lattice point. In CBF the gauge-invariant information about gluon degrees of freedom is contained in correlation functions D, D_1 - for bilocal correlators, and similarly for higher correlators. In principle D, D_1 etc should be found from some equations, which result from the minimum of the free energy of the vacuum. At this stage D and D_1 are introduced by hand in the same way as the gluonic condensate (1) has been introduced into the QCD sum rules.

The main problem now is to write physical amplitudes in terms of vacuum correlators (2). To this end we exploit the Schwinger-Feynman technic of path integrals in the proper time. For the quark Green function S(x,y) we write [9]

$$S(x,y) = (m-\gamma_\mu D_\mu)\ G_q(x,y),\ D_\mu = \partial_\mu - igA_\mu \tag{3}$$

where G_q satisfies

$$(m^2 - D_\mu^2 - g\Sigma F)\ G_q(x,y) = \delta(x,y) \tag{4}$$

and

$$(\Sigma F) = \begin{pmatrix} \sigma_i B_i & \sigma_i E_i \\ \sigma_i E_i & \sigma_i B_i \end{pmatrix}. \tag{5}$$

The Green function G_q can be expanded in powers of spin terms (ΣF) (expansion for heavy quarks [10]) or else written in a closed form (for light quarks [11])

$$G_q = G_s + G_s g(\Sigma F) G_s + \dots \tag{6}$$

where G_s is the Green function of a scalar particle in the external (vacuum) field. Here we for simplicity shall expose the CBF method for scalar quarks, and only in the last section we give final formulas for spin effects of light and heavy quarks.

Consider now the Green function of the quark-antiquark system in the vacuum background field. It may be written as an expansion in number of quark loops, the leading term (no additional quark loops) being

$$G(x,y) = \langle G_q(x,y)\ G_q(y,x)\rangle \tag{7}$$

where the average is taken with respect to vacuum background fields. Using the Schwinger-Feynman method G can be rewritten as a path integral, which in case of nonrelativistic quarks is [9]

$$G(x,y) = \int Dz\, D\bar{z}\, \exp(-K)\, \langle W(Q)\rangle . \qquad (8)$$

Here

$$K = \frac{m_1}{2}\int_y^x \dot{z}^2 d\tau + \frac{m_2}{2}\int_y^x \dot{\bar{z}}^2\, d\tau \qquad (9)$$

and the paths of quark $z_i(\tau)$ and antiquark $\bar{z}_i(\tau)$ start at the same point x_i for $\tau = 0$ and end up at y_i for $\tau = T$ (we are in the Euclidean time). Note, that all vacuum field dependence is contained in $\langle W(Q)\rangle$, which is the Wilson loop average

$$\langle W(Q)\rangle = \langle P \exp ig \int_Q A_\mu\, dz_\mu\rangle \qquad (10)$$

and the closed contour Q is formed by paths of quark and antiquark, which are to be integrated upon in (8).

One can see that the $q\bar{q}$ dynamics is defined totally by the Wilson loop (10), so that one has to investigate its properties more carefully. We shall do it for simplicity in an example of Abelian fields; for non-Abelian fields we shall quote only the final result. We shall show that $\langle W(Q)\rangle$ can be expressed in terms of correlators D, D_1 etc (2), so that the $q\bar{q}$ dynamics is defined by them. To this end we use the Stokes theorem

$$\langle W(Q)\rangle = \langle P \exp ig \int_S F_{\mu\nu}(z)\, d\sigma_{\mu\nu}(z)\rangle \qquad (11)$$

and apply to (11) the so-called cluster expansion [12-13, 9] which allows to express the average $\langle W(Q)\rangle$ through the averages of $\langle F(1)...(F(m)\rangle$:

$$\langle W(Q)\rangle = \exp \sum_{m=L}^{\infty} \frac{(ig)^m}{m!} \int_S d\sigma(1)...\int_S d\sigma(m) \langle\langle F(1)...F(m)\rangle\rangle \qquad (12)$$

We have suppressed subscripts in (12); the cumulants introduced in (12) are defined as

$$\langle\langle F(1)\; F(2)\rangle\rangle = \langle\langle 12\rangle\rangle = \langle 12\rangle - \langle 1\rangle\langle 2\rangle \qquad (13)$$

$$\langle\langle 123\rangle\rangle = \langle 123\rangle - \langle 12\rangle\langle 3\rangle - \langle 13\rangle\langle 2\rangle - \langle 1\rangle\langle 23\rangle + 2\langle 1\rangle\langle 2\rangle\langle 3\rangle \qquad (14)$$

Since for homogeneous vacuum $\langle F\rangle = 0$ we can write the lowest order contribution as [14]

$$\langle W(Q)\rangle = \exp\left[-\frac{g^2}{2}\right] \int_S d\sigma_{\rho\mu}(u) \int_S d\sigma_{\sigma\nu}(u')\langle F_{\rho\mu}(u)\, F_{\sigma\nu}(u')\rangle . \qquad (15)$$

The surface S in (15) and (11) can be chosen as the minimal surface of a given contour Q, since $\langle W(Q)\rangle$ does not depend on the shape of the surface (each term of the cluster expansion (12) however may depend on that, for a discussion see [15, 16]).

The correlator in (15) is the same as in(2)., and in case of non-Abelian fields it is convenient to write in the form [9,14]

$$\langle F_{\rho\mu}(u)\Phi(u,u')F_{\sigma\nu}(u')\Phi(u,u')\rangle = \hat{1}\cdot\beta\Big\{(\delta_{\rho\sigma}\delta_{\mu\nu}-\delta_{\rho\nu}\delta_{\mu\sigma})D(u-u') +$$

$$+\frac{1}{2}\left[\frac{\partial}{\partial u_\rho}(h_\sigma\delta_{\mu\nu}-h_\nu\delta_{\mu\sigma}) + \mu\nu \leftrightarrow \rho\sigma\right] D_1(u-u')\Big\}. \qquad (16)$$

Here we introduced $\Phi(u',u) = P\exp\int_{u'}^{u} ig\, A_\mu dz_\mu$, $h = u - u'$, and $\beta = \frac{\langle trF^2(0)\rangle}{12N_c(D(o)+D_1(0))}$.

5. POTENTIAL vs NONPOTENTIAL DYNAMICS

Voloshin [17] (see also [18]) has argued long ago that very heavy quarks $q\bar{q}$ do not interact in the vacuum gluonic fields in a potential-like manner. Later on it was remarked by Marquard and Dosch [12] that the situation depends on the relation between the correlation time of the vacuum T_g and that of quarks T_q. The full derivation of potential and nonpotential cases was done [10,19] and we quote here the corresponding results.

Let T_q be the period of quark orbit motion, and T_g - the correlation distance in D, D_1 (so that D, D_1 depend on $\frac{u-u'}{T_g}$). Then two limiting cases occur:

1) $T_q \gg T_g$ - potential case. Dynamics in the $q\bar{q}$ system is given by a local potential $\varepsilon(R)$ (R - distance between q and $\bar{q}$) [10]

$$\varepsilon(R) = \beta\left\{2r\int_0^r d\lambda\int_0^\infty d\nu\; D(\lambda,\nu) + \int_0^r \lambda d\lambda\int_0^\infty d\gamma[-2D(\lambda,\nu) + D_1(\lambda,\nu)]\right\}. \quad (17)$$

At large R, $\varepsilon(R)$ has a simple linear asymptotics with the leading term independent of D_1:

$$\varepsilon(R) = \sigma R - C_0,\; R \gg T_g \quad (18)$$

where the string tension is

$$\sigma = \frac{\beta}{2}\, d^2 = \frac{\beta}{2}\int_{-\infty}^{\infty} d\lambda \int_{-\infty}^{\infty} d\nu D(\lambda,\nu) \quad (19)$$

and

$$C_0 = \beta\int_0^\infty \lambda d\lambda\int_0^\infty d\nu[2D(\lambda,\nu) - D_1(\lambda,\nu)]. \quad (20).$$

Note that the form (18) coincides with that of Cornell potential [20] which has been introduced on phenomenological grounds.

So D ensures linear confinement, and as one can argue [14, 15], D cannot have perturbative contributions - in the Abelian case. The structure of D_1, on the other hand, is present both for perturbative and nonperturbative contributions.

To take into account perturbative contributions we must add to D_1 in (17) a term D_1(pert), which has at small distances the following form

$$\beta D_1^{pert}(u) = \frac{4\alpha_s}{\pi}\, C_2 (u_\mu^2)^{-2},\; u_\mu^2 \to 0 \quad (21)$$

where C_2 is the quadratic Casimir operator which is equal to 4/3 for fundamental q, $\bar{q}$ charges $N_c = 3$, and in this case we obtain at small distances

$$\varepsilon(R) \approx \beta R^2\int_0^\infty d\nu(D(0,\nu) + \frac{1}{2}\, D\; (0,\nu)\;) - \frac{4\alpha_s}{3R}\,. \quad (22).$$

It is interesting to compare the contribution of bilocal correlators with that of higher order cumulants. From (12) again in the limit $T_q \gg T_g$ one finds the nonperturbative contribution of the form:

$$\varepsilon_m(R) \approx \begin{cases} \sigma_m R - C_m,\; R \to \infty & (23) \\ R^m \text{const},\; R \to 0,\; R \ll T_g & (24) \end{cases}$$

where m refers to the order of cumulant «F(1)...F(m)». One can see that all terms of the cluster expansion contribute to the string tension,

$$\sigma = \sum_{m=2}^{\infty} \sigma_m, \tag{25}$$

while at small distances the bilocal correlator yields the leading contribution.

2) $T_q < T_g$.

In this case (assuming also that $R \ll T_g$) one should insert in (8) (see [14,10]

$$\langle W(Q)\rangle = \exp\left(-\frac{g^2 S^2 \langle trF^2(0)\rangle}{24N_c}\right) \tag{26}$$

where the surface S(Q) can be written in the form

$$S^2 = \int_0^T r_i(\tau)d\tau \int_0^T r_i(\tau')d\tau'. \tag{27}$$

Dynamics in (8) taking into account (26)-(27) is nonlocal in time (bilocal in case of (27) and multilocal for the contribution of higher cumulants). Therefore it is of nonpotential character. In this case one can solve the path integral (8), or solve an equivalent integral equation, which has been written in [20]. Historically this is a typical case where the QCD sum rules [8] are applicable. Finally there is the perturbation theory of Voloshin-Leutwyler [17,18] where one treats vacuum background field as a perturbation.

Let us consider concrete physical systems and compare T_g and T_q. The latter depends on the dynamics. For heavy quarkonia where Coulomb is most important, one can use for T_q the classical Coulomb period,

$$T_q^{coul} \approx \frac{4\pi n^3}{m_q\left(\frac{4}{3}\alpha_s\right)^2} \tag{28}$$

and for lighter quarks (u,d,s) one can use a classical estimate for linear potential (angular momentum L=0)

$$T_q^{lin} \approx 2(m_q)^{1/2}\frac{(M-2m_q)^{1/2}}{\sigma}. \qquad (29)$$

Both estimates yield $T_q > 1$ fm for charmonium and bottomonium, the same result occurs for light quarks ($m_q \approx 0.3$ GeV) from (29). Excited states have even larger T_q.

Now for T_g we have only a preliminary estimate from lattice calculations [21], $T_g \sim 0.17$ fm. We expect that T_g could be in the interval 0.1 fm - 0.5 fm. In this case all known quark-antiquark systems belong to the potential dynamics which probably is one of the reasons of the success of QPM.

6. LIGHT QUARK DYNAMICS

The previous consideration is valid in nonrelativistic case, i.e. for heavy quarks. The case of light quarks has its own specifics. First, one should take into account relativistic dynamics. Second, the chiral symmetry breaking (CSB) is important at least in some channels. We treat in this Section mostly the first problem, following the method of [22].

Using the proper time formalism in the framework of the path integral, we obtain instead of (8)

$$G(x,y) = \int_0^\infty ds \int_0^\infty d\bar{s}\, Dz\, D\bar{z}\, \exp(-K_{rel})\langle W(Q)\rangle \qquad (30)$$

where

$$K_{rel} = m_1^2 s + m_2^2\bar{s} + \frac{1}{4}\int_0^s \dot{z}_\mu^2(\sigma)d\sigma + \frac{1}{4}\int_0^{\bar{s}} \dot{\bar{z}}_\mu^2(\sigma)d\sigma. \qquad (31)$$

The q and $\bar{q}$ trajectories are described by $z_\mu(\sigma)$, $\bar{z}_\mu(\sigma)$ with boundary conditions $z_\mu(0) = \bar{z}_\mu(0) = y_\mu$, $z_\mu(s) = \bar{z}_\mu(\bar{s}) = x_\mu$.

Assuming that the size (time) of the system $R_q(T_q)$ is much larger than T_g (potential case), we replace $\langle W(Q)\rangle$ by its asymptotic value:

$$W(Q)\rangle = \exp(-\sigma_0 S(Q)) \qquad (32)$$

where S(Q) is the minimal surface of the closed contour Q, formed by paths $z_\mu(\sigma)$, $\bar{z}_\mu(\sigma)$. It is convenient for heavy states to treat the center of mass coordinate

$$R_\mu = \frac{s\bar{s}}{s+\bar{s}} \left[\frac{1}{s} z_\mu + \frac{1}{\bar{s}} \bar{z}_\mu\right] \tag{33}$$

as a slowly moving, and take $R_4 \equiv \tau$ as a time parameter instead of σ, $\bar{\sigma}$. Finally, we introduce instead of s, $\bar{s}$ the "moving masses" μ_1, μ_2

$$\frac{s}{T} = \frac{1}{2\mu_1}, \quad \frac{\bar{s}}{T} = \frac{1}{2\mu_2} . \tag{34}$$

Then for large T the integrand of (30) takes the form exp(-B) where the "action" B is

$$B = \int_0^T d\tau \left[\frac{m_1^2}{2\mu_1} + \frac{m_2^2}{2\mu_2} + \frac{\mu_1+\mu_2}{2} + \varepsilon(\tilde{\mu})\right] \tag{35}$$

and $\varepsilon(\tilde{\mu})$ with $\tilde{\mu} = \frac{\mu_1\mu_2}{\mu_1+\mu_2}$ is an eigenvalue

$$H = \varepsilon(\tilde{\mu}) \ ; \ H = -\frac{1}{2\tilde{\mu}} \frac{\partial^2}{\partial u_i^2} + \sigma_0 (u_i u_i)^{1/2} . \tag{36}$$

Therefore the mass of the bound $q\bar{q}$ system is given by

$$M(\mu_1,\mu_2) = \frac{m_1^2}{2\mu_1} + \frac{m_2^2}{2\mu_2} + \frac{\mu_1+\mu_2}{2} + \varepsilon(\tilde{\mu}) \tag{37}$$

where using the steepest descent method, μ_1 and μ_2 are to be found from the equation

$$\frac{\partial M}{\partial \mu_1} = \frac{\partial M}{\partial \mu_2} = 0 . \tag{38}$$

From (37) and (38) we deduce the following important conclusions:

1) Linear Regge trajectories with the standard slope

$$l = \alpha' M^2, \ \alpha' = \frac{1}{8\sigma_0} . \tag{39}$$

2) Appearance of dynamic (constituent) quark mass which is state-dependent.

For $m_1 = m_2$, $\mu_1 = \mu_2 = \mu_0 = \left[\frac{\sigma_0(n_r+l+1)}{2}\right]^{1/2}$ (40)

For $n_r = l = 0$, $\mu_0 = 0.32$ GeV.

3) The dynamics in (36) looks like nonrelativistic quark model with linear potential, however effective momenta are large, $p \sim \mu_0$ and $\varepsilon(\hat{\mu}) \sim 3\mu_0$, so that the motion is highly relativistic.

The important message here is that the actual dynamics looks like the QPM, however may be outside of its region of validity - "nonrelativistic" equation (36) with relativistic momenta. This may be another argument in favour of QPM. In the next Section we shall see that even spin interactions look like those used in QPM.

Now consider baryons made of light quarks. The Green function can be written as follows [23]

$$G(x,y) = \int d\mu(1)d\mu(2)d\mu(3)\langle W_3(x,y)\rangle \tag{41}$$

where

$$d\mu(i) = \int_0^\infty ds_i e^{-m_i^2 s_i} Dz(i) \exp\left[-\frac{1}{4}\int_0^{s_i} (\dot{z}_\mu^{(i)})^2 d\sigma\right] \tag{42}$$

and $\langle W_3(x,y)\rangle$ is the Wilson loop average with contours along quark paths from y to x. Assuming again that the size of baryon R_q, $T_q \gg T_g$, we can use the asymptotics

$$\langle W_3(x,y)\rangle = \exp(-\sigma_0(S_1 + S_2 + S_3)), \tag{43}$$

where S_i is the minimal surface bounded by the path of the i-th quark and the trajectory of the string junction - $z_\mu^{(0)}(\sigma)$. The latter is determined by the requirement that the sum $S_1 + S_2 + S_3$ be minimal for given trajectories.

Two different configurations have been considered. For a symmetric one (quark distances to the string junction are of the same order) the slope of baryon Regge trajectory is $\frac{1}{12\sigma_0}$, while in the asymmetric case, when one quark is much farther away, the slope is the same as for meson trajectory, $\alpha'_B = 1/8\sigma_0$. In this latter case the baryon has lower mass and the structure of the quark-diquark type. The total baryon mass is

$$M = 4\mu^0 + \frac{m_3^2}{2\mu^0} + \frac{m_1^2+m_2^2}{\mu_0} + \varepsilon_\eta \qquad (44)$$

where μ_0 is the same as for meson (40), and ε_η is the binding energy of the diquark.

Several properties of baryons have thus been found [23]:

1) linear Regge trajectories with the same standard slope $1/8\sigma_0$, as for mesons - in agreement with experiment;

2) appearance of dynamic (constituent) mass of diquark and of the third quark, both are almost equal, $\mu_3 \approx \mu_0 \approx \mu_1 + \mu_2$, $\mu(\text{diquark}) \approx \mu_1 + \mu_2 + \varepsilon_\eta$. The masses are state dependent and grow with angular momentum (see Eq.(40));

3) again the dynamics looks like nonrelativistic QPM, but with string junction and not with pairwise qq potentials.

7. SPIN INTERACTIONS

Spin dependent potentials are usually introduced via the nonrelativistic expansion in inverse powers of the quark mass [24]. The most systematic treatment has been done in CBF method in [10] based on the fact that spin operators enter as local insertions on the quark path, see Eq.(6), and in addition linearly in the operators $\gamma_\mu D_\mu$ in (3). With the help of the cluster expansion (12) one can calculate spin-dependent terms also for light quarks. In the potential case when $T_q \gg T_g$, the final result is reasonably simple and can be cast into the familiar form . Consider the $q\bar{q}$ case with quarks of light or heavy masses. We quote the final result [11]. The Green function can be written as

$$G = [m_1 + (\mu_1 + \pi_4)\gamma_4 + \gamma_i\pi_i]\,[m_2 + (\mu_2 - \pi_4\gamma_4 - \gamma_i\pi_i)\,g(x,y,T)] \qquad (45)$$

where g satisfies:

$$\frac{\partial g(x,y,T)}{\partial T} = -Hg, \qquad (46)$$

π_μ is the relative momentum of q, $\bar{q}$ and H is

$$H = [M(\mu,\mu_2) + V_{SD}]\cdot\hat{1}\cdot\hat{1} + H'. \qquad (47)$$

The spin-dependent interaction V_{SD} has the same structure as in nonrelativistic expansion for heavy quarks and we keep notations of Eichten and Feinberg [24]:

$$V_{SD} = \left(\frac{\sigma_i^{(1)}L_i^{(1)}}{4\mu_1^2} - \frac{\sigma_i^{(2)}L_i^{(2)}}{4\mu_2^2}\right)\left(\frac{2}{r}\frac{dV_1}{dr}\right) + \frac{(\sigma_i^{(2)}L_i^{(1)} - \sigma_i^{(1)}L_i^{(2)})}{2\mu_1\mu_2}\frac{1}{r}\frac{dV_2}{dr} +$$

$$+ \frac{\sigma_i^{(1)}\sigma_i^{(2)}}{12\mu_1\mu_2}V_4(r) + \frac{1}{12\mu_1\mu_2}(3\sigma_i^{(1)}n_i\sigma_k^{(2)}n_k - \sigma_i^{(1)}\sigma_i^{(2)})\,V_3(r). \qquad (48)$$

All potentials are expressed through correlators D, D_1 as follows:

$$\frac{1}{r}\frac{dV_1}{dr} = -\beta\int_{-\infty}^{\infty}d\nu\int_0^r\frac{d\lambda}{r}\left(1-\frac{\lambda}{r}\right)D(\lambda,\nu) \qquad (49)$$

$$\frac{1}{r}\frac{dV_2}{dr} = \beta\int_{-\infty}^{\infty}d\nu\int_0^r\frac{\lambda d\lambda}{r^2}\left[D(\lambda,\nu) + D_1(\lambda,\nu) + \lambda^2\frac{\partial D_1}{\partial\lambda^2}\right] \qquad (50)$$

$$V_3 = -\beta\int_{-\infty}^{\infty}d\nu\, r^2\,\frac{\partial D_1(r,\nu)}{\partial r^2} \qquad (51)$$

$$V_4 = \beta\int_{-\infty}^{\infty}d\nu\left[3D(r,\nu) + 3D_1(r,\nu) + 2r^2\frac{\partial D_1}{\partial r^2}\right]. \qquad (52)$$

Several comments are in order:

1) the term V_{SD} enters the chiral structure $\hat{1}\ \hat{1}$; there are other structures, like $\hat{1}\cdot\hat{\gamma}_5$, $\hat{\gamma}_5\cdot\hat{\gamma}_5$ in H', which are numerically less important even for light quarks;

2) the "moving masses" μ_1, μ_2 should be taken at the minimum of the whole Hamiltonian H (47). For heavy quark masses this amounts to replacement $\mu_i \to m_i$. For light quarks there appears again (now spin-dependent) dynamic (constituent) quark mass of the order of 300 MeV;

3) V_{SD} has the same form as that obtained by nonrelativistic expansion. However the form (47) is exact and should not be treated

perturbatively, in contrast to the former derivations;

4) if one inserts in (50-52) $D_1 \to D_1 + D_1^{pert}$ as in (21) one obtains usual perturbative contributions to spin-spin and spin-orbit potentials;

The forms (49-52) are in agreement with lattice calculations [25].

8. CONCLUSIONS

The approach presented here allows to bridge the gap between QCD and hadron models. The resulting picture for quarkonia, mesons and baryons is close to QPM and flux-tube model, but with important modifications.In particular we can now better understand why QPM is so successful in description of completely different hadrons, such as heavy quarkonia and light baryons [3,5]. But the CBF method has much wider scope. It allows in principle to describe all hadronic systems and interaction between hadrons. All that is to be done in terms of vacuum correlation functions - universal characteristics of the vacuum.

On the way to nuclei and, first of all, to the NN interaction, there is still an important property of the QCD vacuum, which should be included in CBF - the spontaneous chiral symmetry breaking - and connected low-energy chiral physics. This question will be discussed in a subsequent publication.

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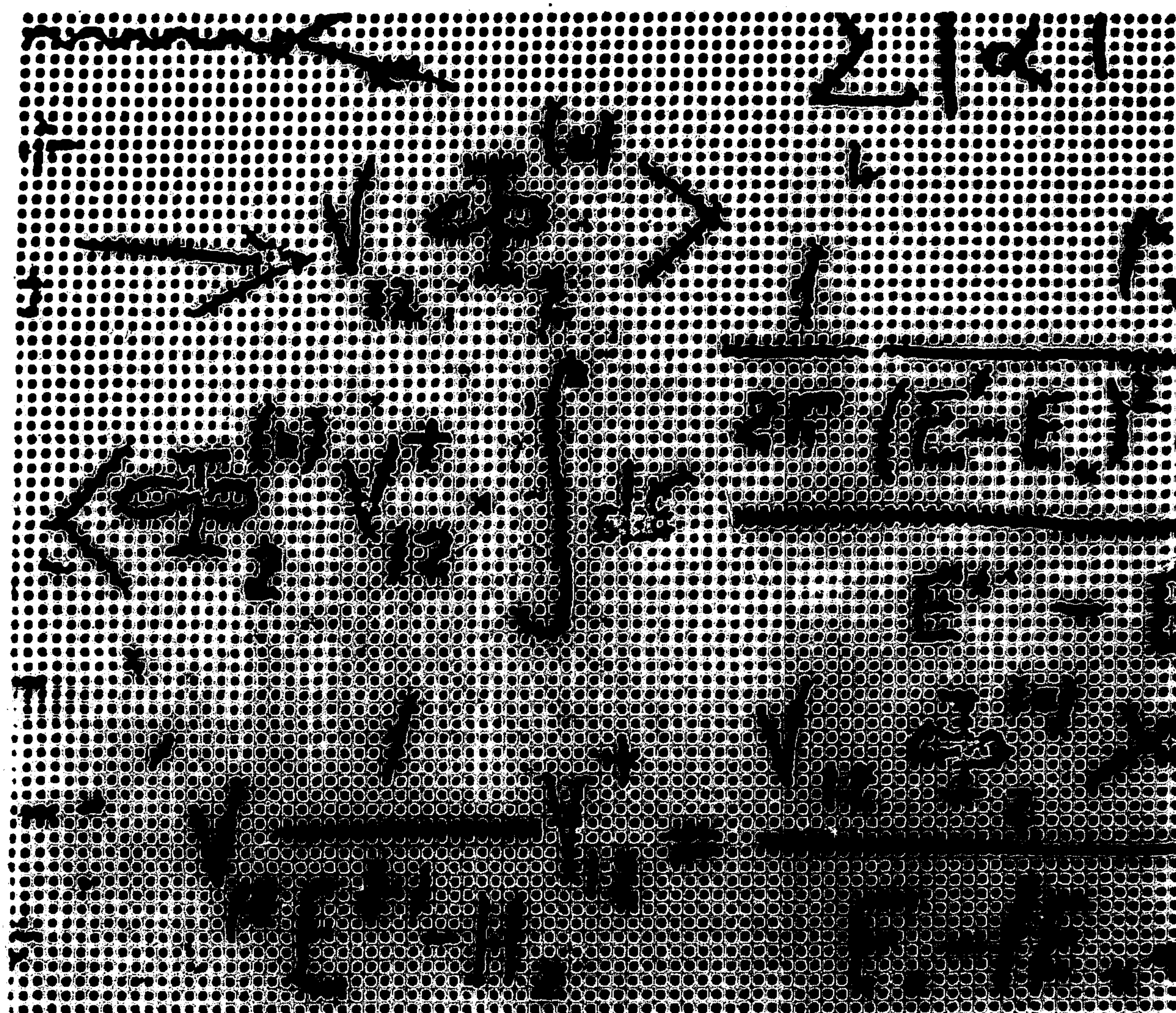

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