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EFFECTIVE INTERACTIONS OF QUARKS AND STATIC PROPERTIES OF HADRONS

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1. Introduction

The potential models of quark interactions were used since the beginning of the quark theory of hadrons[1-3] up to the present time (see, e.g., Refs.[4,5]). In particular, application of the nonrelativistic potential formalism to the heavy-quark bound states ("quarkonia") turned out to be very successful[6]. Extension of the potential picture to the light (u, d, s) - quark dynamics presents an alternative to the picture of the quasifree motion of relativistic quarks in certain bounded region of space, which is known to underlie the bag models[7-9]. To get more insight into the dynamics of a composite relativistic systems, an adequate description of the multiquark (primarily, three - quark) systems may be crucial. Having in mind the difficulties in dealing with the relativistic multiparticle systems as well as uncertainties in our present knowledge of quark potentials, we believe the variational approach to be, at this stage, the most suitable and economic tool in investigation of the dependence of various hadronic observables on free parameters of theory. For this reason, we employ the variational method for a unified approximate treatment of both the two-particle and multiparticle bound systems. For mesons (i.e., the $q\bar{q}$ - states), the numerical solutions, whenever available, of the exact relativistic two-body equations (e.g., the quasi-potential equations[10,11]) may be used as an accuracy test of the approximate variational calculations. For baryons and exotic multiquark states, the nonrelativistic quark model is fruitfully used in line with the relativistic, "one-particle-type" bag models. Success of nonrelativistic models in predicting many relations between the hadronic observables looks rather paradoxically, because the velocities of the constituent quarks have certainly to be relativistic. This work aims to offer, besides simple and, hopefully, effective calculational approach, some arguments underlying the phenomenological utility of nonrelativistic models. We shall not assume the velocities of constituent particles to be small. Yet, the energy functional constructed in Section 3 has apparently quasinonrelativistic form and by redefinition of appropriate parameters may be identified with the well-known nonrelativistic expression. This feature of the proposed scheme enables us also to attain the favourable balance between a comparatively simple treatment of the multiparticle aspects of the problem, analogously to nonrelativistic approach, and consideration of motion of constituent quarks relativistically, the feature used to be dealt with in the framework of the independent particle models.

2. Quark Interaction Potentials

A) Two-body $q\bar{q}$ - potential

To perform calculations, it is first necessary to fix the Lorentz - Dirac properties and radial dependence of the $q\bar{q}$ - interaction potential. The potential should ensure the quark confinement and be consistent, as far as possible, with the known or expected properties and peculiarities of the fundamental field theory - quantum chromodynamics (QCD). According to QCD short distances are dominated by the one-gluon exchange, i.e., as $Q^2 \to \infty$ the vector potential in the momentum representation is

$$V_v(Q^2) \sim \left[Q^2 \ln \frac{Q^2}{\Lambda^2}\right]^{-1}, \tag{2.1}$$

where Λ is some unknown dimensional parameter. Richardson/12/ has proposed an ansatz for extrapolation of (2.1) into the region of small Q^2 (i.e., long distances) by means of the change $\ln \vec{q}^{\,2}/\Lambda^2 \to \ln(1+\vec{q}^{\,2}/\Lambda^2)$. As a result, in the coordinate representation the potential has the asymptotic behaviour $\sim (r \ln r)^{-1}$ and $\sim r$ at short $(r \to 0)$ and long $(r \to \infty)$ distances, respectively, and at the same fixed Λ well describes properties of J/ψ and Υ -mesons. To see at which values of r the asymptotic regime sets in, it is convenient to consider the analytic approximation

$$V_v(r) = \begin{cases} a_v \cdot r, & r \geq r_2 \\ b \cdot \ln \frac{r}{r_0}, & r_1 \leq r \leq r_2 \\ c \cdot \left(r \ln \frac{r}{r_0}\right)^{-1}, & r \leq r_1 . \end{cases} \tag{2.2}$$

The conditions of the continuity of $V(r)$ and $V'(r)$ at points $r = r_{1,2}$ reduce the number of free parameters in (2.2) to one. Constant c is assumed to be known from QCD: the one-loop approximation to the gluon propagator gives

$$c = 8\pi/(33-2n_f), \tag{2.3}$$

where n_f is the number of "effective" quarks, i.e., those with mass $m_q \ll 1/r_1$.
Other constants in (2.2) can be expressed, e.g., in terms of a_v

$$r_0 = e^2 r_1 = \frac{r_2}{e} = \frac{1}{2}\left(\frac{ec}{a_v}\right)^{1/2}$$
$$b = \frac{1}{2}\left(e^3 a_v c\right)^{1/2}, \tag{2.4}$$

where e = 2,718... is the base of natural logarithms. In a region of reasonable values of a_v corresponding to Λ in (2.1) from the range $0,1 \leq \Lambda \leq 0,5$ GeV the values of r_1 and r_2 vary within limits $0,22 \geqslant r_1 \geqslant 0,44$ fm and $4,4 \geqslant r_2 \geqslant 0,88$ fm, and therefore the wave functions of most hadrons should be expected to concentrate mainly in the range of action of the logarithmic potential. Further, considering that the approximation $\ln r/r_0 \simeq \frac{1}{2}\left(r/r_0 - r_0/r\right)$ is well enough in the range $r_1 \leq r \leq r_2$ and has a qualitatively correct asymptotic behaviour for $r > r_2$ and $r < r_1$, we take the expression

$$V_v(r) = \frac{1}{2} b \left(\frac{r}{r_0} - \frac{r_0}{r}\right) \equiv a r - \frac{\varkappa}{r} \tag{2.5}$$

as an approximate representation of $V_v(r)$ at all r. The constant $\varkappa \simeq c$ is independent of a_v and at $n_f = 2$ $\varkappa \simeq 0.8$. In addition to the vector potential, $q\bar{q}$ - interaction should involve also terms with different Lorentz - Dirac properties. Starting from pioneering works/1-3/ a possible role of the scalar interaction potential is repeatedly discussed in current literature. It seems worth noting that the evidence for insufficiency of only one vector potential comes from an old theoretical work/15/ reporting the absence of discrete spectrum of eigenvalues of the Dirac equation with pure-vector potentials of confinement. The Dirac equation may be used for the approximate description of properties of two-particle systems composed of one "very heavy" quark (e.g., b - quark with mass $m_b \simeq 5$ GeV) and one light quark ($m_q \simeq 0$). The very existence of such mesons together with assumed validity of the potential approach testifies to the existence of the universal scalar potential of confinement. From QCD it should be expected that the role of the scalar potential will diminish, as

$r \to 0$, compared to the vector potential. For simplicity, we take the simple linear radial dependence of $V_s(r)$ at all r :

$$V_s(r) = a_s \cdot r. \tag{2.6}$$

Analysis[16] of the spin-orbital splitting of $^3P_{0,1,2}$ -levels of charmonium agrees with the approximate equality $a_s \simeq a_v$. Thus, formulae (2.5) and (2.6) fix our "working" $q\bar{q}$ - potential which we consider universal, i.e., independent of the flavour quantum numbers like isospin, hypercharge, charm, etc.

B) Effective interaction of quarks in many-body systems

According to QCD and confinement hypothesis, all the hadrons are singlets, (anti) quarks are (anti) triplets and gluons are octets of the colour group $SU(3)_c$. The one-gluon exchange potential of a quark two-body interaction contains the colour factor ($\vec{\lambda}_i \cdot \vec{\lambda}_j$) that fixes the strength of the Coulomb interaction of an i-th and j-th quark in an arbitrary, colour-singlet system relatively the Coulomb $q\bar{q}$ - interaction in the meson. Relation of this kind between the phenomenological $q\bar{q}$-potential and quark interaction in many-body systems is, in general, unknown. As a possible approach to solve this problem we resort to the self-consistent field method. Let us consider an N - quark system, being a colour singlet by definition, as a quasi-two-body system composed of a colour point-like triplet (i.e., quark) and spatially distributed colour antitriplet, consisting of (N - 1) quarks. The radial dependence of interaction of the quark with the volume element of the effective "antiquark" is assumed to be known according to the hypothesized independence of the $q\bar{q}$ - interaction of particle flavours (that is, the baryon charge, isospin, etc.). The expectation value of the potential in a given state of the whole system is expressed via the same parameters we introduce to determine the wave function of the N - particle system, and define, in the variational approach, through minimizing the energy of system. Consider now this scheme as applied to the 3-quark state (baryon). The potential V(1,2,3) dependent on quark configuration variables is represented by

$$V(1,2,3) = K \cdot \left[V(1;2,3) + (1 \leftrightarrow 2) + (1 \leftrightarrow 3) \right], \tag{2.7}$$

where V(i; j, k) is the interaction potential of an i-th quark with a spatially extended "antiquark" formed of j-th and k-th quarks and factor K is included to get rid of possible multiple counting of the same terms (for instance, K = 1/2 if each term V(i; j, k) is reduced to the sum of two-body potentials V(i, j) and V(i, k)).

Each term in the sum (2.7), e.g., V (3; 1,2) has the form

$$V(3;1,2)=\int d^3r'\,\rho(r')\,V_0(|\vec{r}_3-\vec{r}\,'|), \tag{2.8}$$

where $\vec{r}\,'$ is the radius - vector of the volume element of the colour antitriplet composed of 1st and 2nd quarks, V_0 stands for our "fundamental" $q\bar{q}$ - potential fixed in the preceding section.

We assume that the distribution function $\rho(r)$ of the colour charge is spherically symmetric, obeys the unit normalization condition and has the form

$$\rho(r)=\frac{1}{2}\int\prod_i d^3r'\,|\Psi(\vec{r}_1\,',\vec{r}_2\,',\vec{r}_3\,')|^2\cdot\left[\delta(\vec{r}-\vec{r}_1\,')+\delta(\vec{r}-\vec{r}_2\,')\right], \tag{2.9}$$

where $\Psi(\vec{r}_1,\vec{r}_2,\vec{r}_3)$ is the symmetric ground-state wave function of the 3-quark system.

Consider the expectation value of V(3; 1,2) in the ground state of baryon

$$\langle V(3;1,2)\rangle=\int|\Psi(\vec{\rho}_3\,',\vec{\lambda}_3\,')|^2\cdot\left[V_0\left(\left|\sqrt{\tfrac{3}{2}}\,\vec{\lambda}_3+\tfrac{1}{\sqrt{2}}\vec{\rho}_3\,'\right|\right)+\right.$$

$$\left.+V_0\left(\left|\sqrt{\tfrac{3}{2}}\,\vec{\lambda}_3-\tfrac{1}{\sqrt{2}}\vec{\rho}_3\,'\right|\right)\right]\cdot|\Psi(\vec{\rho}_3,\vec{\lambda}_3)|^2\,d\vec{\rho}_3\,d\vec{\lambda}_3\,d\vec{\rho}_3\,'\,d\vec{\lambda}_3\,', \tag{2.10}$$

where we introduced the Jacobi coordinates

$$\begin{aligned}\vec{\rho}_3&=\frac{1}{\sqrt{2}}(\vec{r}_1-\vec{r}_2)\equiv\frac{1}{\sqrt{2}}\vec{r}_{12}\\ \vec{\lambda}_3&=\frac{1}{\sqrt{6}}(\vec{r}_1+\vec{r}_2-2\vec{r}_3)\end{aligned} \tag{2.11}$$

and performed integration eliminating the δ - functions.

As a first approximation we take now a test wave function in the factorized form

$$\Psi(\vec{\rho}_3,\vec{\lambda}_3)=\Phi(\vec{\rho}_3)\,\chi(\vec{\lambda}_3)\quad, \tag{2.12}$$

the factorization being fulfilled for two other possible sets of the Jacobi coordinates ($\vec{\rho}_1,\vec{\lambda}_1$) and ($\vec{\rho}_2,\vec{\lambda}_2$) resulting from (2.11) by the change of particle indices. In this case one can obtain

$$\langle V(3;1,2)\rangle=\frac{1}{2}\left\langle\left[V_0(\vec{r}_{13})+V_0(\vec{r}_{23})\right]\right\rangle \tag{2.13}$$

without specifying the form of functions Φ and χ in (2.12). Allowing for the factor $K = 1/2$ in (2.7) we find that on the class of test functions satisfying the condition (2.12) the complete potential can be written in the form

$$V(1,2,3) = \sum_{j>i=1}^{3} V_0(\vec{r}_{ij}) . \tag{2.14}$$

Generalization to the N - body case is straightforward: replacing the normalization factor $1/2$ in the formula of type (2.9) by $1/N-1$, with the same considerations we get

$$V(1,2,\dots N) = \frac{1}{N-1} \sum_{j>i=1}^{N} V_0(\vec{r}_{ij}) \tag{2.15}$$

under the condition of factorization of $\Psi(\vec{\rho}_1,\dots \vec{\rho}_{N-1})$ in either of the Jacobi coordinates

$$\vec{\rho}_n = [n(n+1)]^{-1/2} \cdot \left[\sum_{i=1}^{n} \vec{r}_i - n\vec{r}_{n+1} \right], \quad 1 \le n \le N-1 . \tag{2.16}$$

The potential (2.15) can be used at the first step of the iteration process. Next steps will give rise to the effective many-body forces.

Note one more important point. The presence of δ - functions in (2.9) is equivalent to the assumption that the colour charge in the baryon is carried out only by point-like quarks. Assume now that the instantaneous distribution of the colour charge in the effective antitriplet represents a "string", i.e., the charge is uniformly, for simplicity, distributed along the straight line between points 1 and 2. In this case, even in the zeroth approximation, that is on the class of factorizable test functions, we have

$$\langle V(3;1,2) \rangle = \frac{1}{2} \left\langle \int_{-1}^{+1} ds\, V_0\left(\left| \sqrt{\tfrac{3}{2}}\, \vec{\lambda}_3 - \frac{s}{\sqrt{2}}\, \vec{\rho}_3 \right| \right) \right\rangle , \tag{2.17}$$

where scalar parameter S varies in the range $-1 \le S \le +1$ and defines the coordinate of a point on the string, $\vec{r}(s) = \frac{1}{2}(\vec{r}_1 + \vec{r}_2 - s\,\vec{r}_{12})$. Integration in (2.17) leads in general to the complicated 3-body potentials.

In the following we confine ourselves, while carrying out calculations, to the simplest case of two-body potentials (2.15) of effective qq - and $q\bar{q}$ - interaction in many-body systems.

3. Variational Approach to the Description of Relativistic Systems

Let H be the Hamiltonian of a system of N particles with a two-body interaction. Define operators $\hat{h}_i$ by the equality

$$H=\sum_i T_i+\sum_{i<j}V(i,j)\equiv\sum_i\hat{h}_i=\sum_i\left[T_i+\frac{1}{2}\sum_{j(\neq i)}V(i,j)\right], \tag{3.1}$$

where T(V) are operators of the kinetic (potential) energy. We consider the standard variational problem of finding the minimum of the energy functional

$$\langle\Psi|H|\Psi\rangle=\sum_{i=1}^{N}\langle\Psi|\hat{h}_i|\Psi\rangle\rightarrow extr \tag{3.2}$$

on a given class of normalizable test functions Ψ . Let ε_i be a one-particle energy which corresponds to the expectation value of $\hat{h}_i$ and has the meaning of the Lagrange multiplier. From the obvious inequality

$$\langle\Psi|(\hat{h}_i-\varepsilon_i)^2|\Psi\rangle\geqslant 0 \tag{3.3}$$

we have

$$\langle\Psi|\hat{h}_i|\Psi\rangle\leqslant\langle\Psi|\frac{1}{2\varepsilon_i}(\varepsilon_i^2+\hat{h}_i^2)|\Psi\rangle . \tag{3.4}$$

Now we make a physical assumption on the approximate stationarity of the one-particle energy. If this is really the case, then the left-hand side of inequality (3.4) is close to the right-hand side and instead of the initial energy functional (3.2) we can consider the expression that is close to and majorizes it

$$\sum_{i=1}^{N}\langle\Psi|\frac{1}{2\varepsilon_i}(\varepsilon_i^2+\hat{h}_i^2)|\Psi\rangle\rightarrow extr . \tag{3.5}$$

Let us now specify the form of $\hat{h}_i$ allowing for the relativistic nature of the system. For $\hat{h}_i^2$ the expression is postulated which results from squaring of the Dirac operator for a given particle moving in the external field of N-1 fixed centers with the coordinates of particles of the system and subsequent transition to the two-component formalism. To illustrate this statement, we write down the one-particle Dirac equation with scalar (V_s) and vector (V_v) potentials

$$(\vec{\alpha}\cdot\vec{p}+\beta m+\beta V_s+V_v-\varepsilon)\psi=0 . \tag{3.6}$$

Squaring it we get

$$(\vec{p}^{\,2}+(m+V_s)^2-V_v^2+2\varepsilon V_v-i(\vec{\alpha}\cdot\vec{\nabla})V_v+i\beta(\vec{\alpha}\cdot\vec{\nabla})V_s)\psi=\varepsilon^2\psi . \tag{3.7}$$

The extra components of the wave function can be eliminated and the two-component formalism is introduced via the definition of the chiral components $\Psi_{R(L)}$ and two-component functions φ and χ

$$\Psi = \Psi_L + \Psi_R = \frac{1}{2}(1+\gamma_5)\Psi + \frac{1}{2}(1-\gamma_5)\Psi \tag{3.8}$$

$$\Psi_R = \begin{pmatrix} \varphi \\ -\varphi \end{pmatrix} \qquad \Psi_L = \begin{pmatrix} \chi \\ \chi \end{pmatrix}. \tag{3.9}$$

As far as the spin-orbital term in (3.7) originated by the scalar potential V_s does not commute with γ_5, thus preventing decoupling of Ψ_R and Ψ_L, we restrict ourselves to the "zeroth" approximation in all spin effects and identify $\hat{h}_i^2$ with the operator in the left-hand side of the equation (3.7), where all spin terms are omitted. This is then the operator of the Klein-Gordon-Fock one-particle equation. With the formula (2.15) we obtain the energy functional of the N-particle (quark and antiquark) system

$$E = \langle \Psi_{\vec{P}=0} | \sum_{i=1}^{N} \frac{1}{2\varepsilon_i} \Big\{ \varepsilon_i^2 + \vec{p}_i^{\,2} + \Big[m_i + \frac{1}{2(N-1)} \sum_{j(\neq i)} V_s(i,j)\Big]^2 - \\ - \Big[\frac{1}{2(N-1)} \sum_{j(\neq i)} V_v(i,j)\Big]^2 \Big\} + \frac{1}{N-1} \sum_{i<j} V_v(i,j) \; | \Psi_{\vec{P}=0} \rangle, \tag{3.10}$$

where $\Psi_{\vec{P}=0}(\vec{\rho}_1, \ldots \vec{\rho}_{N-1}; \{\alpha\}; \{\varepsilon_i\})$ are test wave functions dependent on $(N-1)$ Jacobi coordinates, with total momentum of all particles $\vec{P} = 0$, $\{\alpha\}$ is a set of the variable parameters. Besides $\{\alpha\}$, ε_i are also varied to minimize the energy.

The conditions of extremum

$$\frac{\partial E}{\partial \alpha_\ell} = 0 \quad , \quad \ell = 1, 2, \ldots \tag{3.11}$$

$$\frac{\partial E}{\partial \varepsilon_i} = 0 \quad , \quad i = 1, 2, \ldots N \tag{3.12}$$

compose a system of equations for defining the variational parameters. In the nonrelativistic limit, $v_i/c \ll 1$, the expression (3.10) transforms into the known expression of nonrelativistic theory

$$E_{NR} = \langle \Psi | \sum_{i=1}^{N} (m_i + \vec{p}_i^{\,2}/2m_i) + \sum_{i<j} V(i,j) | \Psi \rangle. \tag{3.13}$$

Note, that if in (3.10) we merely redefine the quantity ε_i and call it the "effective mass" whose value is defined by the empirical fit, then we come to an expression, formally equivalent to (3.13) but with a more complex potential function. This observation allows us to get a further insight into the nature of the striking adequacy of many relations between observables, which stem from the application to hadrons of the nonrelativistic quark model (see, in this connection, also Ref. /17/).

4. Calculation of Static Characterictics of Hadrons

A) Hadron Masses

The hadron masses are calculated by minimizing the functional (3.10) with

$$V_s(i,j) = a\, r_{ij} \tag{4.1}$$

$$V_v(i,j) = a\, r_{ij} - \frac{\varkappa}{r_{ij}}$$

$\varkappa = 0.8$ and a is a free parameter.

The relativistic squaring of the potentials (4.1) might suggest the choice of test wave functions with a singularity at $r_{ij} \to 0$ and also a possible redifinition of the normalization condition. Reserving these possibilities we, however, consider, as a first approximation, the simplified version

$$\Psi_{N=2} = E_2 \cdot \exp\left(-\tfrac{1}{2}\gamma r\right) \tag{4.2}$$

for the $q\bar{q}$ - system and

$$\Psi_{N\geq 3} = E_N \cdot \exp\left[-\tfrac{1}{2}\alpha^2\left(\vec{\rho}_1^{\,2} + \ldots + \vec{\rho}_{N-1}^{\,2}\right)\right] \tag{4.3}$$

for systems with the total number of quarks and antiquarks $N \geq 3$, where normalization constants E_N are defined in a usual (i.e., nonrelativistic) way.

The model results should be compared with definite combinations of masses of physical hadrons with the contributions of the spin- dependent potentials eliminated. Let us introduce

$$M = M^{(0)} + \langle V_\sigma \rangle \tag{4.4}$$

$$V_\sigma(1,2,\ldots N) = \sum_{i>j} \mu_i^c \mu_j^c \,(\vec{\sigma}_i \cdot \vec{\sigma}_j)\, V_\sigma(i,j), \tag{4.5}$$

where $M^{(0)}$ is a hadron mass with the "switched-off" spin-spin potential V_σ, μ_i^c is the "colour magneton" of an i-th quark, characterizing the spin-dependent constant at the vector interaction vertex.

According to (4.4), for the meson ground states we have

$$P^{(0)} = V^{(0)} = \frac{3V + P}{4}, \tag{4.6}$$

where particle symbols denote their masses.

For baryons, consisting of two identical $(q_1 = q_2 = q_\ell)$ and one unlike ($q_3 = q_u$) quark, we have the general expression

$$\langle V_\sigma \rangle_B = L(B)\cdot\gamma_{\ell\ell}(B) + U(B)\cdot\gamma_{\ell u}(B), \tag{4.7}$$

where (unspecified) $\gamma_{ij}(B)$ are defined by averaging (4.5) over the radial wave function, and the readily calculable numerical coefficients L(B) and U(B)-over the spin-unitary part of the baryon wave function, corresponding to an approximate SU ($2n_f$) symmetry. If we further put

$$\gamma_{\ell\ell} = \begin{cases} \gamma_{qq}, & \ell = q \\ \gamma_{ss}, & \ell = s \end{cases} \tag{4.8a}$$

$$\gamma_{\ell u} = \gamma_{qs}, \quad (q \equiv u,d),$$

where γ's are no longer dependent on the baryon type in the 56-plet of SU(6) symmetry, then with the factorization

$$\gamma_{qq}\cdot\gamma_{ss} = \gamma_{qs}^2 \tag{4.8b}$$

assumed by (4.5), we get

$$N^{(0)} = \frac{1}{2}(N + \Delta) \tag{4.9a}$$

$$\Lambda^{(0)} = \Sigma^{(0)} = \Lambda + \frac{1}{2}(\Delta - N) = \frac{1}{2}(\Sigma^* + \Sigma) - \frac{1}{6}(\Delta - N - \Sigma^* + \Sigma) = \\ = \frac{1}{2}(\Sigma^* + \Sigma) - \frac{1}{4}(\Sigma - \Lambda) \tag{4.9b}$$

$$\Xi^{(0)} = \frac{1}{2}(\Xi^* + \Xi) + \frac{1}{6(\Delta - N)}\cdot(\Sigma^* - \Sigma)\cdot(\Delta - N - \Sigma^* + \Sigma) \tag{4.9c}$$

$$\Omega^{(0)} = \Omega - \frac{1}{2(\Delta-N)}\cdot(\Sigma^{*}-\Sigma)^{2} \tag{4.9d}$$

$$\Lambda_c^{(0)} = \Sigma_c^{(0)} = \Lambda_c + \frac{1}{2}(\Delta-N), \tag{4.9e}$$

where $\Lambda_c(qqc)$ is the charmed baryon with isospin $I = 0$. Calculated values E and "experimental" (i.e., given by Eqn.'s (4.9a-e)) masses $M^{(0)}$ of mesons and baryons are collected in Table 1. Masses of hadrons in (4.9a-e) were taken from Ref./18/. Underlined entries were used to define free parameters: a = 0.055GeV^2 m_s = 0.33 GeV, m_c = 1.65 GeV. The masses of u- and d-quarks were put zero. For the mass of b-quark we, following/19/ assumed value $m_b = m_c + 3.45 = 5.1$ GeV. Table 1 presents also values of γ and α characterizing the spatial extension of wave functions and values of ε_i and $\varepsilon_{\ell(u)}$. In baryons, ε_ℓ refers to the like quarks and ε_u to the unlike quark. Except for the mass of $\Upsilon(b\bar{b})$ the model predictions are in good agreement with

Table 1

Masses of the $q\bar{q}$ - and q^3 states and parameters γ, ε_1, ε_2 for mesons, and α, ε_ℓ, ε_u for baryons. All values are in units of GeV

Particle	$\gamma(\alpha)$	$\varepsilon_1(\varepsilon_\ell)$	$\varepsilon_2(\varepsilon_u)$	E_{th}	$M^{(0)}_{exp}$
$\pi(q_1\bar{q}_2)$	0.703	0.325	0.325	0.605	<u>0.611</u>
$K(q_1\bar{s}_2)$	0.915	0.405	0.577	0.796	<u>0.793</u>
$\varphi(s_1\bar{s}_2)$	1.109	0.624	0.624	0.953	$\leq$1.02
$D_c(q_1\bar{c}_2)$	1.359	0.579	1.758	1.962	<u>1.97</u>
$F_c(s_1\bar{c}_2)$	1.626	0.783	1.833	2.067	2.10
$J/\psi(c_1\bar{c}_2)$	2.615	2.0	2.0	3.02	3.06
$D_b(q_1\bar{b}_2)$	1.908	0.801	5.205	5.33	-
$F_b(c_1\bar{b}_2)$	2.26	1.01	5.22	5.40	-
$\Upsilon(b_1\bar{b}_2)$	7.02	5.87	5.87	8.98	$\leq$9.46
$N(q^3)$	0.298	0.319	0.319	1.11	1.088
$\Sigma(q^2s)$	0.344	0.356	0.511	1.28	1.27
$\Xi(s^2q)$	0.387	0.53	0.39	1.42	1.44
$\Omega(s^3)$	0.428	0.55	0.55	1.55	1.61
$\Sigma_c(q^2c)$	0.419	0.42	1.753	2.505	2.43
$\Sigma_b(q^2b)$	0.45	0.447	5.176	5.932	-

experiment. For the $b\bar{b}$ - system the wave function is mainly concentrated in the region of transition of the logarithmic potential into the modified Coulomb potential, and therefore, the accepted approximation (2.5) fails to work. The model allows simple relations to be derived for the change of hadron masses with a small change of masses of constituent quarks.

Let us estimate the average value of u - and d - quark masses with the help of the relation (see, e.g., ref. [20])

$$\Sigma_{\pi N}(0) = E_N(m_q) - E_N(0) = m_q \left.\frac{\partial E_N(m_q)}{\partial m_q}\right|_{m_q \to 0} \qquad (4.10)$$

between σ - term $\Sigma_{\pi N}(0)$ in the πN scattering amplitude and the nucleon-mass shift due to a small change of $m_q = \frac{1}{2}(m_u + m_d)$ from zero to a certain finite value. The explicit form of the energy functional (3.10) for nucleons obtained with the use of Eqs. (4.1) and (4.3), enables us to get

$$\left.\frac{\partial E_N}{\partial m_q}\right|_{m_q \to 0} = 3\sqrt{\frac{2}{\pi}}\,\frac{a}{\varepsilon_N \alpha_N} \simeq 1.38 \qquad (4.11)$$

If we now assume [20] $\Sigma_{\pi N}(0) = 51 \pm 5$ MeV, then $m_q \simeq 40$ MeV follows from (4.10) and (4.11).

Consider now the mass splitting of particles, belonging to the same isotopic multiplet, caused by the isotopic mass-difference of d- and u- quarks. Denote by $\delta P = \delta V = E(d\bar{Q}) - E(u\bar{Q})$ the meson-mass shift (the baryon-mass shift δB being analogously defined) due to the nonzero mass difference $\Delta m_{ud} = m_d - m_u \neq 0$. Quantities δP and δB depend on flavours of particles (strangeness, charm,etc.). The ratios of δP and δB for different particles are proportional to those of the quantities $1/\varepsilon_q \gamma$ and $1/\varepsilon_q \alpha$ (cf. Equation (4.11)) which can be found in Table 1

$$\delta(\rho^0\omega) : \delta K : \delta D_c : \delta D_b = 1 : 0.62 : 0.29 : 0.15 \qquad (4.12)$$

$$\delta N : \delta\Sigma : \delta\Xi : \delta\Sigma_c : \delta\Sigma_b = 1 : 0.78 : 0.63 : 0.54 : 0.47. \qquad (4.13)$$

From (4.12) it follows, for instance, that the mass difference of K^0- and K^+- mesons, caused by more heavy mass of the d-quark as compared to that of the u-quark, is twice that of $D_c^+(\bar{d}c)$ and $D_c^0(\bar{u}c)$ caused by the same origin. The quantity $\delta(\rho^0\omega)$ in (4.12) stands for the contribution of the internal u - d mass difference to the nondiagonal mass of the $\rho^{\circ}\omega$ - transition. It is interesting to note that despite a considerable breaking of SU(3) symmetry $(\delta N \neq \delta\Sigma \neq \delta\Xi)$ differences between δB are almost completely compensated in the Coleman - Glashow relation

$$p - n + \Xi^{\circ} - \Xi^{-} - \Sigma^{-} + \Sigma^{+} = 0 \tag{4.14}$$

derived earlier within exact SU(3).

B) Magnetic Moments and M1-Transitions

External electromagnetic fields are introduced in the same way as in the two-component theory of Dirac particles. The vector potential is included through the usual gauge-invariant substitution in the kinetic term of the energy operator (3.10)

$$\vec{p}_i^{\,2} \equiv (\vec{\sigma}_i \cdot \vec{p}_i)(\vec{\sigma}_i \cdot \vec{p}_i) \rightarrow \left[\vec{\sigma}_i \cdot (\vec{p}_i + e_i \vec{A}(x_i))\right]\left[\vec{\sigma}_i \cdot (\vec{p}_i + e_i \vec{A}(x_i))\right] \tag{4.15}$$

In an external, uniform magnetic field there appears the correction to energy linear in field strength $\mathcal{H}$:

$$\delta E(\mathcal{H}) = -\left\langle \sum_i \frac{e_i}{2\varepsilon_i} (\vec{\sigma}_i \cdot \vec{\mathcal{H}}) \right\rangle \equiv -\vec{\mu} \cdot \vec{\mathcal{H}} \, . \tag{4.16}$$

From (4.16) it is seen that instead of the i-th quark mass the effective quark magneton contains ε_i . Averaging the magnetic moment operator over spin wave functions and applying the numerical values of ε_i from Table 1 we obtain the magnetic moments of baryons. Formulae for magnetic moments and the transition moments $\mu_{B^*B} = \langle B(1/2^+) | \hat{\mu} | B^*(3/2^+) \rangle$ coincide in form with formulae of the nonrelativistic quark model /21/ with the only difference: the quark magnetons are defined not by the universal parameters (masses of quarks) but by the quantities ε_i which are different for a given sort of quark residing inside different hadrons, i.e., depend on the quark environment. From Table 2, where for comparison we report also results of the nonrelativistic quark model /21/, it is seen that the relativistic consideration makes better agreement of theory with experiment /18,22/. Consider now relations for decay rates of the magnetic dipole (M1) transitions $V \rightarrow P\gamma$. Transitions of this kind between "light" hadrons are characterized by a great energy release, and therefore, the calculation of matrix elements requires to take into account the relativistic contraction of the wave functions of moving particles, the retardation effects, etc. We take a simplified scheme of calculation, assuming that in the ratios of radiative widths unknown dynamical form factors dependent on the photon energy will approximately cancel out and a remaining factor will be proportional to the ratio of the transition moments calculated for the "unphysical" situation when the initial and final mesons are at rest. Then we can calculate the $V \to P\gamma$ decay widths by the following formula

$$\Gamma(V \to P\gamma) = \Gamma(\omega \to \pi^0\gamma) \cdot \left[\frac{\mu(VP)}{\mu(\omega\pi^0)} \right]^2 \cdot \left[\frac{m_\omega \cdot (m_V^2 - m_P^2)}{m_V \cdot (m_\omega^2 - m_\pi^2)} \right]^3 . \qquad (4.17)$$

For $\Gamma(\omega \to \pi^0\gamma) = 789 \pm 92$ KeV [22,30] the values calculated by (4.17) are collected in Table 3. The transition magnetic moments $\mu(VP)$ are calculated by using formulae of the nonrelativistic theory and values of ε_i from Table 1. The inclusion of the dependence of μ_q and μ_s on the mass of quarks- neighbours almost eliminates the discrepancy existing till now between theory (see, e.g., Ref [30]) and experiment for the decays $K^* \to K\gamma$ and $\varphi \to \eta\gamma$.

Table 2

Magnetic moments of baryons in units of nuclear magnetons

Baryon	μ_{NR} [21]	μ_{rel}	μ_{exp} [17,22]
p	input	2.94	2.97
n	-1.86	-1.96	-1.91
Λ	input	-0.61	-0.614 ± 0.005
Σ^+	2.67	2.55	2.33 ± 0.13
Σ^-	-1.09	-0.97	-1.41 ± 0.25
$(\Lambda\Sigma^0)$	1.61	1.52	$1.82^{+0.25}_{-0.17}$
Ξ^0	-1.44	-1.32	-1.237 ± 0.016
Ξ^-	-0.49	-0.52	-0.75 ± 0.06
Ω^-	-1.83	-1.71	-

Table 3

Radiative decay widths of vector mesons in units of KeV. The angles of the $\omega\varphi$- and $\eta\eta'$- mixing are $\theta_V = 40°$ and $\theta_P = -10°$, respectively

Decay	Γ_{th}	Γ_{exp} /22,30/
$\omega \to \pi^0 \gamma$	$\underline{789 \pm 92}$	789 ± 92
$\rho^- \to \pi^- \gamma$	86 ± 10	67 ± 7
$K^{*0} \to K^0 \gamma$	87 ± 10	75 ± 35
$K^{*-} \to K^- \gamma$	52 ± 6	62 ± 14
$\varphi \to \pi^0 \gamma$	8.6 ± 1.0	6.5 ± 1.9
$\varphi \to \eta \gamma$	47 ± 6	67 ± 9
$J/\psi \to \eta_c \gamma$	1.68	-
$D_c^{*0} \to D_c^0 \gamma$	14.8	-
$D_c^{*+} \to D^+ \gamma$	0.23	-
$F_c^{*+} \to F_c^+ \gamma$	0.011	-

For the radiative decays of heavy mesons ($J/\psi \to \eta_c \gamma$, $D^* \to D \gamma$, $F^* \to F \gamma$) the energy release is small, the retardation and recoil effects can safely be neglected, and therefore we have calculated the absolute values of widths by the formula of the long-wave approximation

$$\Gamma(V \to P\gamma) = \frac{1}{3\pi} \cdot \left(\mu_{q_1} - \mu_{\bar{q}_2}\right)^2 \cdot \left(\frac{m_V^2 - m_P^2}{2m_V}\right)^3, \qquad (4.18)$$

where $\mu_q = e_q / 2\varepsilon_q$, $m_{V(P)}$ are meson masses.
Radiative decays of heavy mesons manifest in a clear-cut way the specific prediction of the model - a significant change of the effective magnetons of u, d, s - quarks depending on the mass of the spectator quark. A strong, compensation-type decrease of width of the $F^* \to F\gamma$ decay may produce an extremely interesting and unexpected situation when the weak-decay branching ratios of the F^* -meson will no longer be negligible.

C) Quarks Dimensions of Hadrons

Let us first assume that quarks are point-like ("bare") particles and define the charge radius $\langle r^2\rangle^{b}_{ch}$ by

$$\langle r^2\rangle^{b}_{ch} = \langle \sum_{i=1}^{N} e_q(i)\cdot(\vec{r}_i - \vec{R})^2 \rangle$$
$$\vec{R} = \sum_i \varepsilon_i \vec{r}_i \Big/ \sum_i \varepsilon_i \,, \tag{4.19}$$

where $e_q(i)$ is the i-th quark electric charge and averaging runs over the ground-state wave function.
For charged particles the r.h.s. of (4.19) should be divided by the total charge of the system.

As is seen from Table 4, numerical values of $\langle r^2\rangle^{b}_{ch}$ are, on the average, smaller than experimental charge radii of hadrons. The

Table 4

Charge radii $\langle r^2\rangle^{b}_{ch}$ and $\langle r^2\rangle^{exp}_{ch}$ in units of fm^2.

Particle	$\langle r^2\rangle^{b}_{ch}$	$\langle r^2\rangle^{exp}_{ch}$
$\pi^\pm$	0.24	0.43 ± 0.03 /31/
$K^\pm$	0.16	0.28 ± 0.05 /32/
K^0	-0.03	-0.054 ± 0.026 /33/
D_c^+	0.058	-
D_c^0	-0.085	-
p	0.44	0.74 ± 0.02 /34/
n	0	-0.12 ± 0.01 /35/
Σ^+	0.42	-
Σ^-	0.33	-
Ξ^-	0.26	-
Ω^-	0.21	-

way out of the difficulty is rather obvious and consists in supplying the hadron valence quarks with internal structure. As is known, the quark-parton model of deep inelastic interactions exposes the multi-component state vector of the hadron: the hadron consists of the "valence" quarks-partons and of a "sea" of gluons and $q\bar{q}$-pairs. On the other hand, the static properties of hadrons are described under the assumption that hadrons are formed out of a minimal number of constituent quarks. For consistency with the above mentioned multi-component picture of the hadron, it is now necessary to supply the valence quarks themselves with the multi-component state vector. The dynamical degrees of freedom of "sea" quarks and gluons are phenomenologically reflected in properties of potentials of quark interactions and in the existence of the internal structure, non-zero effective mass, and so on, for the valence quarks. It should, furthermore, be expected that in the peripheral region of distribution of the hadronic matter there come into effect the clusterization of the coloured quarks and gluons into "colourless" virtual hadrons. In this sense, one may speak about a "pion cloud" and mesonic degrees of freedom of hadrons /23/. Contributions of meson (primarily, pion) currents into the hadron characteristics should stronger depend on hadron quantum numbers and their masses compared to the intrinsic structure of valence quarks caused by fluctuations of the accompanied gluon field.

The relation

$$\langle r^2\rangle_p + \langle r^2\rangle_n = \langle r^2\rangle_q + \langle r^2\rangle_p^{b} + \langle r^2\rangle_n^{b} \tag{4.20}$$

and the values collected in Table 4 give for the intrinsic quark radius $\langle r^2\rangle_q = 0.18 \pm 0.02\ \mathrm{fm}^2$ (q = u, d).

The isovector pion current gives no contribution into the isoscalar combination (4.20) but should be essential in forming the negative value of the experimental neutron charge radius. If we parametrize the meson radii by the expressions

$$\langle r^2\rangle_{\pi^+} = \langle r^2\rangle_q + \langle r^2\rangle_\pi^{b} \tag{4.21a}$$

$$\langle r^2\rangle_{K^+} = \frac{2}{3}\langle r^2\rangle_q + \frac{1}{3}\langle r^2\rangle_s + \langle r^2\rangle_{K^+}^{b} \tag{4.21b}$$

$$\langle r^2\rangle_{K^0} = -\frac{1}{3}\langle r^2\rangle_q + \frac{1}{3}\langle r^2\rangle_s + \langle r^2\rangle_{K^0}^{b} \tag{4.21c}$$

and make use of the numerical values of $\langle r^2\rangle_{\pi(K)}^{exp}$ and $\langle r^2\rangle_{\pi(K)}^{b}$

in Table 4, then we get $\langle r^2\rangle_q = 0.17 \pm 0.03$ fm^2 and $\langle r^2\rangle_s = 0.07 \pm 0.07$ fm^2. Within large errors, these estimates confirm qualitative expectations.

D) Exotic Multiquark Hardons

Properties of hadrons for which the dominant configuration is the multiquark state are extensively discussed in literature[24-27]. Almost all works are based on the MIT bag model[7]. Here we apply the developed scheme to calculate masses $M_N^{(o)}$ of (not yet unambigously identified) hadrons with total number N of light (u- and d-) quarks and antiquarks being $3 < N \leq 12$ ($N = 12$ is the maximal number of quarks which can be in the ground state with the total orbital angular momentum $L = 0$ and the symmetric radial wave function). Results of our calculation are collected in Table 5 where, by the "triality" principle, the state with $N = 4$ should be understood as $(q^2\bar{q}^2)$, $N = 5$ - as $(q^4\bar{q})$, and so forth. Numerical values of $M_N^{(o)}$ do not take into account the spin-dependent interactions and therefore their direct comparison with calculations including those interactions is difficult. Meanwhile, in comparing all models with the real situation one should take in view a possibly serious inadequacy caused by a too restrictive way of imposing the confinement conditions. Multiquark states admit the formation of "colourless" clusters out of a smaller number of constituenty quarks and for their relative motion the mechanism of

Table 5

Masses of multiquark states and parameters α and ε in units of GeV.

N	4	5	6	9	12
$\alpha(q)$	0.282	0.273	0.268	0.261	0.258
$\varepsilon(q)$	0.323	0.326	0.328	0.330	0.332
$E_N(q)$	1.557	1.996	2.43	3.74	5.04
$\alpha(s)$	0.401	0.387			
$\varepsilon(s)$	0.557	0.560			
$E_N(s)$	2.16	2.75			

confinement should not necessarily operate. However, the test wave function (4.3) we are using does not allow for that fact. Besides, in the bound states containing antiquarks the processes of virtual $q\bar{q}$ - annihilation may lead to a complex mixture of hadron states with the same quantum numbers but different number of constituents. These factors make it not a simple task to reveal a possible existence, in the hadron-hadron scattering channel, of the multiquark quasistationary states and to identify their composite structure. In view of the signature, which facilitate their experimental observation and of relative isolation from other states, permitting a more reliable interpretation, a search for resonances formed only out of strange quarks and antiquarks and decaying into φ - mesons looks especially promising, e.g., $(s^2\bar{s}^2) \to 2\varphi$. The possible evidence for existence of such a state with mass 2.35 GeV comes from the data on the reaction π^-(22 GeV) + p $\to n + 2\varphi$ /28/. In Table 5 we present several values of $M_N^{(0)}$ for hadrons constructed only from $s(\bar{s})$ - quarks. One can notice that the difference between the mass of $(s^2\bar{s}^2)$- state and the doubled mass of $(s\bar{s})$ - state amounts to 0.32 GeV that is close to the "experimental value": 2.35 - 2 x 1.02 = 0.31 GeV, if the possible resonance with mass 2.35 GeV/28/ is indeed interpreted as the $(s^2\bar{s}^2)$ - state. In nuclear processes with large energy-momentum transfers (see, e.g., ref./29/) an important role may be played by virtual many-baryon clusters with the mean density of "quark matter". For the estimation of the probability of formation of such clusters in real nuclei of great inportance is the ratio of the average density of "quark" matter to the usual "nuclear" matter. Let R_N be the radius of sphere with an equivalent uniform distribution of the matter density for an N-quark state

$$R_N^2 = \frac{5}{3}\langle r^2\rangle_N = \frac{5}{2}\cdot\frac{N-1}{N}\cdot\frac{1}{\alpha_N} . \tag{4.22}$$

Then, the wanted ratio is

$$\rho_N/\rho_A = \frac{R_A^3 E_N}{R_N^3 \cdot A m_p} . \tag{4.23}$$

Taking m_p = 0.94 GeV, $R_A = r_0 A^{1/3}$ (r_0 = 1.2 fm) and corresponding values of $E_N \equiv M_N^{(0)}$ and α_N from Table 5, we find that the ratio (4.23) varies from 3.7 to 6 for the quark states with baryon number varying from B = 2 to B = 4. This is considerably larger than corresponding values obtained in the MIT - bag model (there the ratio is almost constant and equals 2).

5. Conclusion

We summarize now the main results of our consideration.

1. The potential picture of the light quark dynamics together with the universality hypothesis for the $q\bar{q}$ - potential gives very encouraging results. The consistent nonperturbative inclusion of the spin-dependent quark interactions is necessary for further development of the scheme.

2. In the multiquark systems the self-consistent field approximation results naturally in the iterative definition of the state-dependent effective quark interactions. In the first approximation the effective two-body potential is derived, which is related through Equation (2.15) with the basic qq - potential.

3. The principal result of the variational approach proposed, Equation (3.10) for the energy functional and the extremum conditions (3.11) and (3.12), unifies the virtues of the nonrelativistic models, consisting in the comparative simplicity of dealing with the multiparticle features of the problem, and the relativistic description of the constituent quark motion. The quasinonrelativistic form of Equation (3.10) offers more insight into the nature of striking adequacy of many relations between the hadron observables derived from nonrelativistic quark models.

4. The model formulated indicates the effects of more deep breaking of the flavour unitary symmetries. We take the mass difference of the constituent quarks to be the only origin of the $SU(n)_{flav.}$ symmetry breaking. The qualitatively new result of the model, which reflects the multiparticle treatment of the problem and seems not to be implied by the independent particle models, is the variation of static parameters of a given quark due to environment, especially by mass of the spectator quarks. For example, the effective magnetons of the light quarks and the isotopic mass splitting of hadrons due to internal mass difference of the u- and d- quarks are predicted to change considerably while we go from light hadrons to heavy ones. These results provide a better agreement of theory and experiment for the magnetic moments of baryons and the radiative widths of the magnetic dipole transitions in light hadrons. The future measurement of the $D^*_{c(b)} \to D_{c(b)}\gamma$, $F^*_{c(b)} \to F_{c(b)}\gamma$ decay widths will be crucial for testing the model predictions because here the effects are especially strong.

5. The "quark" radii of hadrons, dependent on the spatial extension of the quark wave functions, are substantially smaller

than the experimental charge radii of hadrons. This means that the constituent valence quarks are "dressed" particles with their own internal structure.

6. In the model considered the masses of the multiquark states are somewhat larger while the density of the "quark matter" is much larger than the values given by the MIT bag model.

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