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QUARK-DIQUARK APPROXIMATION OF THE THREE-QUARK STRUCTURE OF BARYONS IN THE QUARK CONFINEMENT MODEL

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1. Introduction

The three-quark structure of baryons seems to be most available in the unitary classification of hadrons [1]. However, the three-body problem is very difficult and can be solved only under some simplifications (for example, in nonrelativistic quark models with the simplest potential [2] and bags [3]). But, there is a possibility to reduce the three-body problem to the two-body one. It is sufficient to suggest that one of the constituent particles of a baryon can be regarded as a quark and the other as a tightly bound state of two quarks, a diquark.

In the original work [1] on the quark model of mesons and baryons Gell- Mann discussed the possibility of existence of free diquarks. Later, the quark-diquark model of baryons was suggested by Ida and Kobayashi [4] and independently by Lichtenberg and Tassie [5]. The supersymmetric generalization of the quark-diquark model has been performed by Catto and Gursey in [6]. It founded on a suggestion that the chromomagnetic fields between the quarks and diquarks were the same in the excited mesons and baryons. In the papers [7,8], the diquark approach was applied to describe exotic and charm mesons. Chan and Hogassen [9] treated dibaryon and multibaryon states using diquarks and larger quark clusters. Cahill and collaborators [10] have suggested the QCD-hadronization which allows appearance of diquark variables. The concept of diquarks was also useful in treating deep inelastic scattering [11].

Thus, the quark-diquark model is a quite suitable approximation of the three- quark structure of baryons. Despite this fact, little has been done along the line of actual quantitative calculations of baryon characteristics in a dynamical quark-diquark model. Mainly, all calculations were directed for getting the static characteristics of baryons: masses, magnetic moments, etc. [12,13].

More sophisticated characteristics of baryons are form factors and phase scattering which are defined by their inner structure. One should know the hadronization mechanism and quark behaviour at large distances for the description of these values.

We have developed [14]-[17] the quark confinement model (QCM) based on the definite representations about the hadronization and quark confinement. First, hadrons are treated as collective colourless excitations of quark-gluon interactions. Second, the quark confinement is realized as averaging over the vacuum gluon backgrounds. Strong, weak and electromagnetic hadron interactions can be described in the QCM from a unique point of view. The preliminary calculations [14]-[17] of the meson and baryon processes have shown that the model reproduces the quark structure of hadrons quite accurately. The hadron inner structure in the QCM is defined by the quark behaviour at large distances.

We have considered [15] a nucleon and a Δ-isobar composed of three quarks and calculated the electromagnetic and strong meson-baryon form factors. The results were in agreement with experiment and other approaches.

At the same time, the consideration of baryons as three-quark systems encounters some difficulties caused by the nature of two-loop diagrams defining the baryon form factors. These diagrams are the convolution of the entire functions [14]-[16] which leads to the growth of physical matrix elements at high energies. The special assumptions were used to avoid this problem. It has been remarked that the S-matrix can be constructed without the above-mentioned difficulties when there are only three lines at the vertex of the interaction Lagrangian (for example, meson+quark+antiquark). Therefore, the hypothesis arised to describing a baryon as a quark-diquark system. In this case, we have only three lines of the vertex of the interaction Lagrangian (baryon+quark+antiquark).

In this paper, we consider the possibility of quark-diquark approximation of the three-quark structure of baryons in the framework of the QCM. The idea of such an approximation is based on the following physical picture. We start with the three-quark structure of octet and decuplet of baryons. The $SU(3)$-quark currents with baryon quantum numbers are symmetric with respect to the permutation of quark fields. Here, only one quark effectively takes part in an interaction of baryons with other fields as mesons, leptons, photons. Therefore, the other two quarks can be considered as a hard cluster, a diquark. Thus, a baryon can be represented as a composition of two particles, quark and diquark. Further, this representation will be called the quark-diquark approximation of the three-quark structure of baryons.

In this paper, we will use this approximation for description of the main low-energy characteristics of baryons as magnetic moments, electromagnetic radii and form factors, ratios of axial and vector costants in semileptonic baryon decays, strong form factors and decay widths.

Earlier [17], the NN-scattering phase shifts were calculated using the strong meson-nucleon form factors obtained in this approach.

Our results are in agreement with experimental data and other approaches. Thus, a quark-diquark approximation of the three-quark structure of baryons is available for taking into account an inner structure of baryons at low energies.

The paper is organized in the following way. In Sec.2 the main notions of the QCM are given. In Sec.3 a set of hadron quark currents, used in the calculations of physical matrix elements is deduced. The connection of the QCM with other approaches is illustrated. In Sec.4 a quark-diquark approximation of the three-quark structure of baryons is given in detail. In Sec.5 the numerical results for the main baryon properties are deduced and the comparison with the experimental data and other approaches are performed. The further investigations are discussed. In Appendix the technique of calculating one-loop quark and quark-diquark diagrams in the QCM is given.

2. The Basic Notions of the QCM

The first assumption of the QCM is related with some notion about hadronization. It is assumed that hadron fields arise after integration over gluon and quark variables in the QCD generating functional [18,19]. Let us demonstrate this transition. The generating functional can be written in the form

$$Z_{QCD} = \int \delta q \int \delta \bar{q} \int \delta B \delta[\partial B] \Delta_L[B] exp\{iS[B]\} \tag{1}$$

where $\Delta_L[B]$ is the Faddeev-Popov determinant fixing the Lorentz gauge,

$$S[B] = \int dx L_{QCD}(x), \qquad L_{QCD} = \frac{1}{8g^2} tr F^2_{\mu\nu} + \bar{q}(i\hat{\partial} + \hat{B})q, \tag{2}$$

$$F_{\mu\nu} = \partial_\mu B_\nu - \partial_\nu B_\mu + [B_\nu, B_\mu], \quad B_\mu = B^a_\mu t^a \quad tr(t^a t^b) = 2\delta^{ab}.$$

Here, $t^a (a = 1, ..., N_c^2 - 1)$ are the SU(N)-generators, B^a_ν and q are the gluon and quark fields, respectively. The Lagrangian (2) is invariant under transformations

$$B_\mu \to B_\mu^\omega = \omega B_\mu \omega^{-1} + \partial_\mu \omega \omega^{-1}, \qquad q \to q^\omega = \omega q \ .$$

We suggest that nontrivial gluon vacuum backgrounds should provide the confinement of all colour objects. According to this representation we divide the gluonic fields into vacuum backgrounds $B^a_{\mu vac}(x, \sigma_{vac})$ characterized by a set of parameters $\{\sigma_{vac}\}$ and quantum fluctations $b^a_\mu(x)$. A similar decomposition was performed in the work [20]. The gluonic field $B^a_\mu(x)$ is represented in the form

$$B_\mu(x) = (B_{\mu vac}(x, \sigma_{vac}) + b_\mu(x))^{\omega(x)} .$$

The field b_μ is chosen to be in the Lorentz gauge in the background field

$$D_\mu(B_{vac}) b_\mu = 0 \ ,$$

where $D_\mu(B)$ is the covariant derivative. The unity is inserted into the generating functional in the standard manner:

$$1 = \Phi[B] \int \delta b \int \delta\omega \int d\sigma_{vac} \delta[B - (B_{vac} + b)^\omega] \delta[D_\mu(B_{vac}) b_\mu] \ .$$

Here, the gauge invariant functional $\Phi[B] = \Phi[B^\omega]$ is defined by the given equality.

Performing simple transformations one can obtain

$$Z_{QCD} = \int \delta q \int \delta\bar{q} \int d\sigma_{vac} W[J] exp\{i \int dx \bar{q}(i\hat{\partial} + \hat{B}_{vac}) q\}, \qquad (3)$$

$$\begin{aligned} W[J] \ &= \ \int \delta b \delta[D_\mu(B_{vac}) b_\mu] \Phi[B_{vac} + b] \ * \\ &* \ \ exp\{\frac{i}{8g^2} \int dx tr F^2_{\mu\nu}[B_{vac} + b] + i \int dx b^a_\mu J^a_\mu\} \ , \end{aligned} \qquad (4)$$

where $J^a_\mu = \bar{q}\gamma^\mu t^a q$ is colour quark current. Recalling the definition of the full gluon Green function in the background field

$$G^{a_1 \dots a_n}_{\mu_1 \dots \mu_n}(x_1, \dots, x_n | B_{vac}) = \frac{\delta^n ln W[J]}{\delta J^{a_1}_{\mu_1}(x_1) \dots \delta J^{a_n}_{\mu_n}(x_n)}$$

we have

$$W[J] = exp\{\sum_n \frac{1}{n!}\int dx_1 ... \int dx_n J^{a_1}_{\mu_1}(x_1)...J^{a_n}_{\mu_n}(x_n) \quad (5)$$
$$* \quad G^{a_1...a_n}_{\mu_1...\mu_n}(x_1...x_n|B_{vac})\}.$$

Inserting (5) into (3) one can obtain

$$Z_{QCD} = \int \delta q \int \delta \bar{q} \int d\sigma_{vac} exp\{i \int dx \bar{q}(i\hat{\partial} + \hat{B}_{vac})q + \sum_n L_n\}, \quad (6)$$

$$L_n = \frac{1}{n!}\int dx_1 ... \int dx_n J^{a_1}_{\mu_1}(x_1)...J^{a_n}_{\mu_n}(x_n) G^{a_1...a_n}_{\mu_1...\mu_n}(x_1, ..., x_n|B_{vac}) .$$

It is to be remarked that the representation (6) is completely equivalent to the initial one (1). For further advancement, the vacuum backgrounds B_{vac} and the full connected gluon Green functions $G^{\{a\}}_{\{\mu\}}(x|B_{vac})$ should be specified.

Let us consider the term L_2 in (6) from which the mesonic fields can be extracted. Emergence of mesons is defined by the behaviour of tne two-point gluon Green function $G^{a_1 a_2}_{\mu_1 \mu_2}(x_1, x_2|B_{vac})$. We suggest that

$$G^{a_1 a_2}_{\mu_1 \mu_2}(x_1, x_2|B_{vac}) = i g_{\mu_1 \mu_2} \delta(x_1 - x_2) G_0 \, \delta^{a_1 a_2} . \quad (7)$$

This suggestion underlies both the Nambu and Jona-Lasinio model [18] and similar approaches [19,21]. Using (7) in L_2, we have

$$L_2 = \frac{i}{2} G_0 \sum_a \int dx (J^a_\mu(x))^2. \quad (8)$$

Perfoming the Firtz transformations in (8) and keeping only the leading $1/N_c$-term, one can obtain

$$L_2 = \frac{i}{2} \sum_{f,J=P,V,S,A} \delta_J c_J \int dx J^2_{Jf}(x). \quad (9)$$

Here,

$$J_{Jf} = \bar{q}\lambda_f \Gamma_J q, \qquad \Gamma_J = I, \gamma^\mu, i\gamma^5, \gamma^\mu\gamma^5 \qquad J = S, V, P, A.$$

$$c_P = c_S = G_0, \qquad c_V = c_A = \frac{1}{2}G_0, \qquad \delta_P = \delta_S = -\delta_V = -\delta_A = 1,$$

λ_f are the Gell-Mann matrices of the flavour $SU(3)$-group. Further, let us use the represention

$$exp\{\pm\frac{i}{2}\int dxJ^2\} = \int \delta M exp\{\mp\frac{i}{2}\int dxM^2 + i\int dxJM\}.$$

Inserting this one into (6) and taking into account only the term L_2 we have

$$Z^{(2)}_{QCD} = \int \delta q \int \delta\bar{q} \int d\sigma_{vac} \int \prod_{Jf} \delta M_{Jf} \tag{10}$$

$$*exp\{-\frac{i}{2}\sum_{Jf}\delta_J \int dx M^2_{Jf} + i\int dx\bar{q}(i\hat{\partial} + \hat{B}_{vac} + M)q\},$$

where

$$M(x) = \sum_{Jf}\sqrt{c_J}M_{Jf}(x)\Gamma_J\lambda^f.$$

After integrating over the quark fields one can obtain

$$\begin{aligned} Z^{(2)}_{QCD} &= \int \prod_{Jf}\delta M_{Jf}\int d\sigma_{vac} exp\{-\frac{i}{2}\sum_{Jf}\int dx\delta_J M^2_{Jf}(x) - \qquad (11)\\ &- \sum_n \frac{N_c}{n}\int dx_1 ... \int dx_n tr[M(x_1)S(x_1,x_2|B_{vac})...\\ &\quad ...M(x_n)S(x_n,x_1|B_{vac})]\}. \end{aligned}$$

Here

$$S(x_1,x_2|B_{vac}) = i(i\hat{\partial} + \hat{B}_{vac})^{-1}\delta(x_1 - x_2).$$

Our next assumption consists in that expression (11) can be written in the form

$$
\begin{aligned}
Z^{(2)}_{QCD} &= \int \prod_{Jf} \delta M_{Jf} exp\{-\frac{i}{2}\sum_{Jf}\int dx \delta_J M^2_{Jf}(x) - \\
&- \sum_n \frac{N_c}{n}\int dx_1 ... \int dx_n \int d\sigma_{vac} tr[M(x_1)S(x_1,x_2|B_{vac})... \\
&* M(x_n)S(x_n,x_1|B_{vac})]\}.
\end{aligned} \tag{12}
$$

It is denoted that all the quark loops at low energies can be connected by the hadron fields but not the gluon vacuum ones.

Further, let use give off the terms diagonalized over meson variables M_J from the expression in the sum in (12) (here the flavour indices are omitted):

$$
\begin{aligned}
&- \frac{i}{2}\int\int dx_1 dx_2 \sum_J \delta_J M_J(x_1)[\delta(x_1-x_2) + \\
&+ \frac{N_c c_J}{(2\pi)^2}\Pi_J(x_1-x_2)]M_J(x_2) = \\
&- \frac{i}{2}\sum_J \delta_J \int dp \tilde{M}^*_J(p)[1+\frac{N_c c_J}{(2\pi)^2}\tilde{\Pi}_J(p^2)]\tilde{M}_J(p)\,,
\end{aligned} \tag{13}
$$

where

$$
\Pi_J(x_1-x_2) = i(2\pi)^2 \int d\sigma_{vac} tr[\Gamma_J S(x_1,x_2|B_{vac})...\Gamma_J S(x_n,x_1|B_{vac})]
$$

$$
\tilde{\Pi}_J(p^2) = \int dx e^{ipx}\Pi_J(x)\,. \tag{14}
$$

Further, we represent the operator $\tilde{\Pi}_J(p^2)$ in the form

$$
\tilde{\Pi}_J(p^2) = \tilde{\Pi}^{ren}_J(p^2) + \tilde{\Pi}_J(m^2_J) + \tilde{\Pi}'_J(m^2_J)(p^2-m^2_J)
$$

and require the performance of the condition

$$
1+\frac{N_c c_J}{(2\pi)^2}\tilde{\Pi}_J(m^2_J) = 0. \tag{15}
$$

In fact, equation (15) gives the connection of the meson mass spectra with the parameters characterizing the behaviour of the full gluon Green function at large distances (7) and confinement properties.

Taking into account (15) we pass to the following normalization in (14)

$$M_J \to M_J[-\frac{N_c c_J}{(2\pi)^2}\tilde{\Pi}'_J(m_J^2)]^{-1/2}.$$

We have

$$Z^{(2)}_{QCD} = \int\prod_J \delta M_J exp\{\frac{i}{2}\int dx \sum_{Jf}\delta_J M_{Jf}(\Box - m_J^2)M_{Jf} - \tag{16}$$
$$-\sum\nolimits'_n \frac{i^n}{n}\int dx_1 ...\int dx_n \int d\sigma_{vac} N_c tr[M(x_1)S(x_1,x_2|B_{vac})...$$
$$...M(x_n)S(x_n,x_1|B_{vac})]\},$$

where

$$M(x) = \sum_{Jf} M_{Jf}(x)\Gamma_J\lambda^f[-\frac{N_c}{(2\pi)^2}\tilde{\Pi}'_{Jf}(m_{Jf}^2)]^{-1/2}.$$

The prime in the sum (16) implies the use of Π^{ren} in the quadratic term over M_J. The representation (16) does not contain the value G_0 which defines the behaviour of the gluon Green function near the point $p^2 = 0$. The generating functional (16) defining the meson-meson interaction by means of the quark loops underlies our model. However, it is more convenient for calculations to use another functional which is completely equivalent to (16). Let us show that the functional (16) can be written in the form

$$Z^{(2)}_{QCD} = \int\prod_J \delta M_J exp\{\frac{i}{2}\sum_J \delta_J \int dx M_J(x)(\Box - m_J^2)M_J(x)\} \tag{17}$$
$$* \int d\sigma_{vac}\int \delta q \int \delta\bar{q} exp\{i\int dx \bar{q}(x)(i\hat{\partial} + \hat{B}_{vac})q(x) +$$
$$+ i\sum_J g_{0J}\int dx M_J(x)\bar{q}(x)\Gamma_J q(x)\}$$

if the wave function renormalization constant of meson M_J is equal to zero.

Indeed, let us integrate over the quark fields in (17) using the same assumptions about the measure $d\sigma_{vac}$:

$$Z^{(2)}_{QCD} = \int \prod_J \delta M_J exp\{\frac{i}{2}\sum_J \delta_J \int dx M_J(x)(\Box - m^2_{0J})M_J(x) - \tag{18}$$

$$-\sum_n \frac{i^n}{n}\int dx_1 ... \int dx_n \int d\sigma_{vac} N_c tr[M(x_1)S(x_1, x_2|B_{vac})...$$

$$...M(x_n)S(x_n, x_1|B_{vac})]\}.$$

Here

$$M(x) = \sum_J g_{0J} M_J(x)\Gamma_J \ .$$

One can give off the quadratic term over M_J in an analogous way

$$\begin{aligned} & \frac{i}{2}\int dp \tilde{M}^*_J(p)[p^2 - m^2_{0J} - h_{0J}\tilde{\Pi}_J(p^2)]\tilde{M}_J(p) = \\ = \ & \frac{i}{2}\int dp|\tilde{M}_J(p)|^2\{(p^2 - m^2_J)Z^{-1}_J - h_{0J}\tilde{\Pi}^{ren}_J(p^2)\} \ , \end{aligned}$$

where

$$h_{0J} = \frac{N_c g^2_{0J}}{(2\pi)^2}, \quad m^2_J = m^2_{0J} + h_{0J}\tilde{\Pi}_J(m^2_J), \quad Z^{-1}_J = 1 - h_{0J}\tilde{\Pi}'_J(m^2_J).$$

Performing in (18) the replacement $M_J \to Z^{1/2}_J M_J$ and introducing the renormalized constant $h_J = Z_J h_{0J}$, we get the functional (16) if

$$h_J = Z_J h_{0J} = \frac{h_{0J}}{1 - h_{0J}\tilde{\Pi}'_J(m^2_J)} = \frac{1}{[-\tilde{\Pi}'_J(m^2_J)]}.$$

This equality can be true if $h_{0J} \to \infty$, that is

$$Z_J = 1 + h_J\tilde{\Pi}'_J(m^2_J) = 0. \tag{19}$$

It is the so-called compositeness condition in quantum field theory [22].

Thus, the representation (17) with the condition (19) is completely equivalent to (16) and will underlie our model.

The hadrons can be constructed in the same manner. These states should arise from the terms L_n for $n > 2$ in (6). As in the case L_2, the gluon Green function $G^{\{a\}}_{\{\mu\}}(x|B_{vac})$ is supposed to be analytical at point $p = 0$. The product of colourless n-quark states, which can be identified with the corresponding hadrons, is given off the L_n by using the Firtz transformations. But, this way is very difficult. Therefore, we will use equivalence of the representations (16) and (17) with auxiliary condition (19). We will start from the hadron spectrum with the given masses and quantum numbers. Then, the hadron-quark interaction Lagranqian $L_H(x)$ will be constructed by using the quark composition of hadrons:

$$L_H(x) = g_H H(x) J_H(x), \tag{20}$$

where $J_H(x)$ is the quark current with the quantum numbers of H.

It is more convenient in practice to use the S-matrix instead of the generating functional (17):

$$S = \int d\sigma_{vac} T exp(i \int dx L_H(x)) \, . \tag{21}$$

The time-order product in (21) is the Wick standard T-product for the hadron and quark fields. The quark propagator has the following form:

$$S(x_1, x_2|B_{vac}) =< 0|T(q(x_1)\bar{q}(x_2))|0 >= i(\hat{p} + \hat{B}_{vac})^{-1}\delta(x_1 - x_2). \tag{22}$$

All hadron interactions will be described by the quark diagrams induced by S-matrix (21) averaged over vacuum backgrounds. The coupling constant g_H is determined from the compositeness condition (19).

The following basic assumptions of the QCM are about quark confinement. It is proposed that the averaging over vacuum background fields B_{vac} of the quark diagrams generated by the S-matrix (21) should provide the quark confinement and make the ultraviolet finite theory.

The confinement ansatz in the case of one-loop diagrams describing the meson-meson interactions consists in the following:

$$\int d\sigma_{vac} tr[M(x_1)S(x_1x_2|B_{vac})...M(x_n)S(x_nx_1|B_{vac})] \rightarrow$$
$$\rightarrow \int d\sigma_v tr[M(x_1)S_v(x_1-x_2)...M(x_n)S_v(x_n-x_1)]. \tag{23}$$

Here

$$S_v(x_1-x_2) = \int \frac{d^4p}{(2\pi)^4 i} e^{-ip(x_1-x_2)} \frac{1}{v\Lambda_q - \hat{p}} \ .$$

The parameter Λ_q characterizes the confinement range. The measure $d\sigma_v$ is defined as

$$\int \frac{d\sigma_v}{v-z} = G(z) = a(-z^2) + zb(-z^2) \ . \tag{24}$$

The function $G(z)$ called the confinement function is an entire analytical function which decreases faster than any degree of z in an Euclidean direction $z^2 \rightarrow -\infty$. $G(z)$ is a universal function, i.e., is dependent on colour and flavour. In other words, the function $G(z)$ is unique for all quark diagrams defining the hadron interaction at low energies. The choice of $G(z)$ is one of the model assumptions. However, as calculations have showed, only integral characteristics of the function $G(z)$ are important for the description of low-energy physics [14,16].

The simplest shapes of $a(u)$ and $b(u)$ were used:

$$a(u) = a_0 exp(-u^2 - 2a_1 u) \qquad b(u) = b_0 exp(-u^2 + 2b_1 u) \ . \tag{25}$$

The parameters $\{a_i, b_i\}$ and the dimensional one Λ_q were defined by fitting over well-established experimental data [23]. It was found that the best description of the experimental data was achieved for $a_0 = b_0 = 2$, $a_1 = 0.5$, $b_1 = 0.2$ and $\Lambda_q = 460 \quad Mev$.

3. The Quark Structure of Hadrons in the QCM

Let us adduce the interaction Lagrangians for the mesons as two-quark systems and the baryons as three-quark ones.

1. Mesons:

$$L_M = \frac{g_M}{\sqrt{2}} \sum_{i=1}^{8} M_i \bar{q} \Gamma_M \lambda_i q \ . \tag{26}$$

Here M_i is the Euclidean fields connected with the physical ones in a standard manner [23]; λ_i (i=1,...,8) are the Gell-Mann matrices ($\lambda_0 = \sqrt{2/3I}$); Γ_M are the Dirac matrices:

$i\gamma^5$ for pseudoscalar mesons $P(\pi, K, \eta, \eta')$;

γ^μ for vector mesons $V(\rho, K^*, \omega, \phi)$;

$\gamma^\mu\gamma^5$ for axial mesons $A(a_1, K_1, f_1)$;

$I - iH_S\hat{\partial}/\Lambda_q$ for scalar mesons $S(a_0, K_0, f_0, \varepsilon)$. The necessity of introducing the auxiliary term with a derivative to the scalar quark current was established in detail in [16].

Mixing octet-singlet angles are defined as

$$(\eta', \omega, \varepsilon) \to \cos\delta_\Gamma\left(\frac{\bar{u}u + \bar{d}d}{\sqrt{2}}\right) - (\bar{s}s)\sin\delta_\Gamma$$

$$(\eta, \pi, f_0) \to -\sin\delta_\Gamma\left(\frac{\bar{u}u + \bar{d}d}{\sqrt{2}}\right) - (\bar{s}s)\cos\delta_\Gamma$$

$\delta_\Gamma = \theta_\Gamma - \theta_{I\Gamma}$; $\theta_{I\Gamma} = 35°$ is the ideal mixing angle. The mixing angles of pseudoscalar and vector mesons are chosen to be equal to $\delta_P = 46°$ and $\delta_V = 0°$, respectively. The scalar meson parameters δ_S, H_S and m_ε are supposed to be free.

2. Baryons:

The three-quark currents with baryon quantum numbers must be symmetrical with respect to permutation of all quarks. There exist two independent three-quark currents for a baryon octet $\frac{1}{2}^+$. Therefore, the interaction Lagrangians can be written in the form:

$$L_B = L_{BT} + L_{BV}, \tag{27}$$

$$L_{BI} = g_{BI}\bar{B}J_{BI} + h.c. = ig_{BI}\bar{B}^k_j R_I^{kj;j_1,j_2,j_3} q^{a_1}_{j_1} q^{a_2}_{j_2} q^{a_3}_{j_3} \varepsilon^{a_1a_2a_3} + h.c.$$

Here, $I = T, V$; $j = (\alpha, m)$; a_i, α_i, m_i are the colour, spin, flavour indices, respectively.

$$B^k_j = \begin{pmatrix} \frac{1}{\sqrt{2}}\Sigma^0 + \frac{1}{\sqrt{6}}\Lambda^0 & \Sigma^+ & p \\ \Sigma^- & -\frac{1}{\sqrt{2}}\Sigma^0 + \frac{1}{\sqrt{6}}\Lambda^0 & n \\ -\Xi^- & \Xi^0 & -\frac{2}{\sqrt{6}}\Lambda^0 \end{pmatrix}$$

is the octet baryon matrix.

$$q^a_j = \begin{pmatrix} u^a \\ d^a \\ s^a \end{pmatrix} \quad \text{is a set of quark fields.}$$

The matrices $R_I^{kj;j_1j_2j_3}$ can be written in the following forms:

$$\begin{aligned} R_T^{kj,j_1,j_2,j_3} &= \varepsilon^{km_2n}\{6g^{\alpha\alpha_1}\delta^{mm_1}(C\gamma^5)^{\alpha_2\alpha_3}\delta^{nm_3} + \\ &+ 6(\gamma^5)^{\alpha\alpha_1}\delta^{mm_1}C^{\alpha_2\alpha_3}\delta^{nm_3} - \\ &- \frac{1}{3}(\sigma^{\mu\nu}\gamma^5)^{\alpha\alpha_1}\delta^{mm_1}(C\sigma^{\mu\nu})^{\alpha_2\alpha_3}\delta^{nm_3} + \\ &+ (\sigma^{\mu\nu}\gamma^5)^{\alpha\alpha_1}\lambda_i^{mm_1}(C\sigma^{\mu\nu})^{\alpha_2\alpha_3}\lambda_i^{nm_3}\}, \end{aligned} \tag{28}$$

$$\begin{aligned} R_V^{kj,j_1,j_2,j_3} &= \varepsilon^{km_2n}\{2g^{\alpha\alpha_1}\delta^{mm_1}(C\gamma^5)^{\alpha_2\alpha_3}\delta^{nm_3} - \\ &- 2(\gamma^5)^{\alpha\alpha_1}\delta^{mm_1}C^{\alpha_2\alpha_3}\delta^{nm_3} - \\ &- (\gamma^\mu)^{\alpha\alpha_1}\delta^{mm_1}(C\gamma^\mu\gamma^5)^{\alpha_2\alpha_3}\delta^{nm_3} - \\ &- \frac{1}{3}(\gamma^\mu\gamma^5)^{\alpha\alpha_1}\delta^{mm_1}(C\gamma^\mu)^{\alpha_2\alpha_3}\delta^{nm_3} + \\ &+ (\gamma^\mu\gamma^5)^{\alpha\alpha_1}\lambda_i^{mm_1}(C\gamma^\mu)^{\alpha_2\alpha_3}\lambda_i^{nm_3}\}. \end{aligned} \tag{29}$$

There is the only interaction Lagrangian for the baryon decuplet:

$$L_D = ig_D\overline{D}^{k_1k_2k_3}_{\mu\alpha}\Gamma^\mu_{\alpha\alpha_1\alpha_2\alpha_3}q^{a_1}_{\alpha_1k_1}q^{a_2}_{\alpha_2k_2}q^{a_2}_{\alpha_2k_2}\varepsilon^{a_1a_2a_3} + h.c., \tag{30}$$

where

$$\Gamma^\mu_{\alpha\alpha_1\alpha_2\alpha_3} = g^{\alpha\alpha_1}(C\gamma^\mu)^{\alpha_2\alpha_3} - \frac{i}{2}(\gamma^\nu)^{\alpha\alpha_1}(C\sigma^{\mu\nu})^{\alpha_2\alpha_3};$$

$$D^{111} = \Delta^{++}, \quad D^{112} = \frac{1}{\sqrt{3}}\Delta^+, \quad D^{122} = \frac{1}{\sqrt{3}}\Delta^0, \quad D^{222} = \Delta^-;$$

$$D^{113} = \frac{1}{\sqrt{3}}\Sigma^{*+}, \quad D^{123} = \frac{1}{\sqrt{6}}\Sigma^{*0}, \quad D^{223} = \frac{1}{\sqrt{3}}\Sigma^{*-};$$

$$D^{133} = \frac{1}{\sqrt{3}}\Xi^{*0}, \quad D^{233} = \frac{1}{\sqrt{3}}\Xi^{*-}, \quad D^{333} = \Omega^-.$$

The field $D_\mu(x)$ satisfies the Rarita-Schwinger equation with the auxiliary conditions:

$$\gamma^\mu D_\mu(x) = 0, \qquad \partial^\mu D_\mu(x) = 0 .$$

One has to note that these baryon quark currents coincide with those that have been used in the QCD sum rule calculations [25] if they are written in the isotopic components.

The electroweak interactions are introduced in a standard manner:

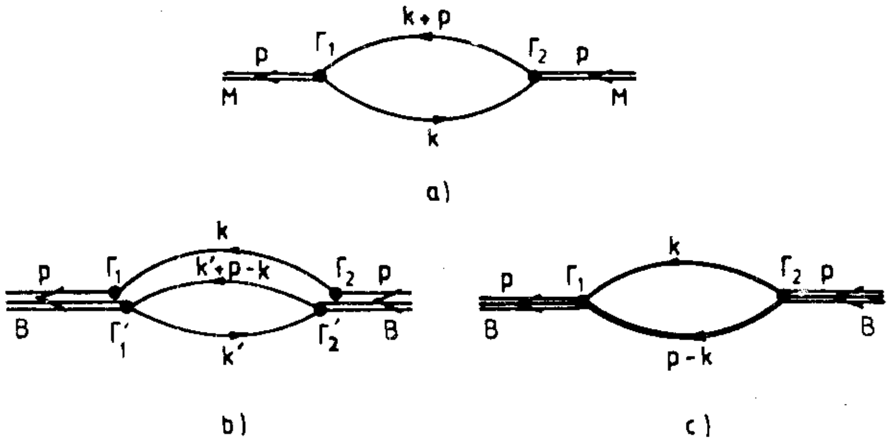
$$L_{em} = eA_\mu \bar{q}_a Q \gamma^\mu q_a, \quad L_{weak} = \frac{G_F}{\sqrt{2}} l_\mu \bar{q}_a J \gamma^\mu (1-\gamma^5) q_a \ , \qquad (31)$$

where A_μ is an electromagnetic field and l_μ is a lepton current.

$Q = diag\{2/3, -1/3, -1/3\}$ is the charge quark matrix;

$$J = \begin{pmatrix} 0 & 0 & 0 \\ \cos\theta_c & 0 & 0 \\ \sin\theta_c & 0 & 0 \end{pmatrix} \quad \text{is the Cabbibo matrix.}$$

The coupling constants $g_H (H = M, B, D)$ are determined from the compositeness condition $Z_H = 1 + 3g_H^2/(4\pi^2)\tilde{\Pi}'_H(m_H) = 0$, where $\tilde{\Pi}'_H$ is the derivative of the hadron mass operator defined by the diagrams, Fig.1a, for mesons and, Fig.1b, for baryons.



1. Mass operators:
   a) two-quark diagram,
   b) three-quark diagram,
   c) quark-diquark diagram.

The interaction Lagrangians (26), (27), (30) and (31) allow one to describe hadron interactions from the unique point of view. One has to emphasize that both the static hadron characteristics as decay widths, magnetic moments etc., and the momentum dependences of physical matrix elements as form factors, phase shifts etc., can be obtained in the QCM.

As has been mentioned above, the confinement function G(z) and dimensional parameter Λ (see, (25)) are free model parameters which have been defined by fitting the well-established experimental data (see, Table 1). One can see there is good agreement with experimental data. It is interesting to consider the limits of zero masses the limits for the values shown in Table 1. One can see that in this case the well-known low-energy relations as the Goldberger-Treiman relation, the ρ-universality hypothesis, the relations between $g_{\pi\gamma\gamma}$ and f_π, $g_{\pi\gamma\gamma}$ and $g_{\omega\pi\gamma}$ are reproduced with an accuracy of 4-5 %.

Table 1
The main low-energy values

Process	Observable value	QCM Fit	QCM Zero masses	Expt. [23]
$\pi \to \mu\nu$	f_π, Mev	132	$3.53\Lambda_q/(4\pi) =$ 129	132
$\rho \to \gamma$	$g_{\rho\gamma}$	0.18	$1.07/(2\pi) =$ 0.17	0.20
$\pi^0 \to \gamma\gamma$	$g_{\pi\gamma\gamma}$, Gev^{-1}	0.260	$0.96\sqrt{2}/(4\pi^2 f_\pi) =$ 0.266	0.276
$\omega \to \pi\gamma$	$g_{\omega\pi\gamma}$, Gev^{-1}	2.10	$2.82\pi g_{\pi\gamma\gamma} =$ 2.36	2.54
$\rho \to \pi\pi$	$g_{\rho\pi\pi}$	6.0	$1/g_{\rho\gamma} = 5.9$	6.1

The auxiliary free parameters H_S, δ_S and ε-meson mass m_ε characterizing scalar mesons in our approach were defined by fitting the experimental data on the $\pi\pi$- and $\pi\gamma$-scattering and scalar meson decay widths [16]. It was found that $H_S = 0.55$, $\delta_S = 17^\circ$ and $m_\varepsilon = 600$ Mev.

In this paper, we will use the same parameters for describing baryon physics.

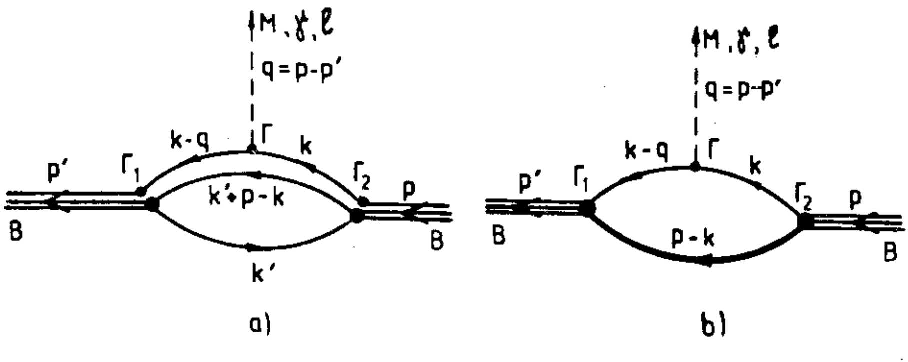
4. Quark-Diquark Approximation of the Three-Quark Structure of the Octet and Decuplet of Baryons

The idea of the quark-diquark approximation of the three-quark structure of baryons is based on the symmetry of the SU(3)-baryon quark currents with respect to the permutation of all quark fields. Effectively, only one quark interacts with other fields due to this symmetry and the other two quarks can be considered as a hard core, a diquark.

This picture is realized in the following way by using the Feynman diagram langauge. The subdiagram corresponding to the independent quark loop

$$\Pi_v^{\Gamma_1\Gamma_2}(p) = \int \frac{d^4k}{4\pi^2 i} tr[\Gamma_1' S_v(p+k)\Gamma_2' S_v(k)] \tag{32}$$

is given off the diagrams Fig.1b and Fig.2a describing mass operators and baryon vertices, respectively, due to such symmetry. Our main assumption consists in that the diagrams in Fig.1b and Fig.2a can be changed to



2. Baryon vertex:
a) three-quark diagram,
b) quark-diquark diagram.

the one-loop quark-diquark diagrams in Fig.1c and Fig.2b, respectively, according to the rule

$$\int \frac{d^4k}{\pi^2 i} \int d\sigma_v \Gamma_1 S_v(k) \Gamma_2 \Pi_v^{\Gamma_1\Gamma_2}(p-k) \Longrightarrow \tag{33}$$
$$\Longrightarrow \int \frac{d^4k}{\pi^2 i} \int d\sigma_v \Gamma_1 S_v(k) \Gamma_2 D_v^{\Gamma_1\Gamma_2}(p-k),$$

$$\int \frac{d^4k}{\pi^2 i} \int d\sigma_v \Gamma_1 S_v(k-q) \Gamma S_v(k) \Gamma_2 \Pi_v^{\Gamma_1\Gamma_2}(p-k) \Longrightarrow$$
$$\Longrightarrow \int \frac{d^4k}{\pi^2 i} \int d\sigma_v \Gamma_1 S_v(k-q) \Gamma S_v(k) \Gamma_2 D_v^{\Gamma_1\Gamma_2}(p-k)\,,$$

where $D_v^{\Gamma_1\Gamma_2}(p)$ is considered to be a propagator of a diquark in the vacuum background fields:

$$D_v^{\Gamma_1\Gamma_2}(k) = \frac{d^{\Gamma_1\Gamma_2}}{v^2\Lambda_D^2 - k^2}\ . \tag{34}$$

Here the additional parameter Λ_D characterizes the diquark confinement region. The parameters $d^{\Gamma_1\Gamma_2}$ are chosen in the forms convenient for calculations:

$$d^{PP} = C_{PP}, \quad d^{SS} = C_{SS}, \quad d^{VV} = C_{VV} g^{\mu\nu}, \quad d^{AA} = C_{AA} g^{\mu\nu},$$

$$d^{TT} = C_{TT} g^{\mu\alpha} g^{\nu\beta}, \quad d^{AP} = -d^{PA} = C_{AP}(ik_\mu),$$

$$d^{VT} = -d^{TV} = C_{VT}(ik_\alpha g^{\mu\beta} - ik_\beta g^{\mu\alpha})\ ,$$

where $C_{\Gamma_1\Gamma_2}$ are numerical coefficients. One has to remark, this quark-diquark approximation can be used for the description of both octet and decuplet of baryons. The general requirement to this approximation consists in that it should not break the Ward identity between the baryon electromagnetic vertex $\Lambda_{\gamma BB}$ and mass operator $\Sigma_B(p)$:

$$\frac{\partial}{\partial p^\mu} \Sigma_B(p) = \Lambda^\mu_{\gamma BB}(p,p)\ .$$

This requirement with taking into account the compositeness condition $Z_B = 0$ gives us the following identity in the case of baryon octet:

$$g_{BT}^2 F_{TT}(p) + g_{BV} g_{BT} F_{VT}(p) + g_{BV}^2 F_{VV}(p) = 0 \tag{35}$$

for any momentum p. This means that the following conditions should be performed simultaneously:

$$\begin{aligned}
F_{TT} &= \int \frac{d^4k}{\pi^2 i} \int d\sigma_v [36 S_v(k) D_v^{PP}(p-k) + 36\gamma^5 S_v(k)\gamma^5 D_v^{SS}(p-k) - \\
&- 3\sigma^{\mu\nu}\gamma^5 S_v(k)\sigma^{\alpha\beta}\gamma^5 D^{TT}_{v;\mu\nu,\alpha\beta}(p-k)] = 0,
\end{aligned}$$

$$\begin{aligned}
F_{VT} &= \int \frac{d^4k}{\pi^2 i} \int d\sigma_v [24 S_v(k) D_v^{PP}(p-k) - 24\gamma^5 S_v(k)\gamma^5 D_v^{SS}(p-k) \\
&+ 6 S_v(k)\gamma^\mu D^{PA}_{v;\mu}(p-k) - 6\gamma^\mu S_v(k) D^{AP}_{v;\mu}(p-k) + \\
&+ 3\sigma^{\alpha\beta}\gamma^5 S_v(k)\gamma^\mu\gamma^5 D^{TV}_{v;\alpha\beta,\mu}(p-k) - \\
&- 3\gamma^\mu\gamma^5 S_v(k)\sigma^{\alpha\beta}\gamma^5 D^{VT}_{v;\mu,\alpha\beta}(p-k)] = 0,
\end{aligned}$$

$$\begin{aligned}
F_{VV} &= \int \frac{d^4k}{\pi^2 i} \int d\sigma_v [4 S_v(k) D_v^{PP}(p-k) + 4\gamma^5 S_v(k)\gamma^5 D_v^{SS}(p-k) + \\
&+ 2 S_v(k)\gamma^\mu D^{PA}_{v;\mu}(p-k) - 2\gamma^\mu S_v(k) D^{AP}_{v;\mu}(p-k) - \\
&- \gamma^\mu S_v(k)\gamma^\nu D^{AA}_{v;\mu\nu}(p-k) + \\
&+ 3\gamma^\mu\gamma^5 S_v(k)\gamma^\nu\gamma^5 D^{VV}_{v;\mu\nu}(p-k)] = 0.
\end{aligned}$$

It turns out that there are no solutions for arbitrary g_{BT} and g_{BV} because the identity $F_{VT} = 0$ cannot be satisfied by any choice of $C_{\Gamma_1\Gamma_2}$. In particular cases, when either $g_{BT} \neq 0$ and $g_{BV} = 0$ or $g_{BT} = 0$ and $g_{BV} \neq 0$, the parameters $C_{\Gamma_1\Gamma_2}$ are defined uniquely

$$C_{\Gamma_1\Gamma_2} = \begin{cases} 1, & \Gamma_1 = \Gamma_2 = S, P, A, \\ 1/3, & \Gamma_1 = \Gamma_2 = V, \\ 2, & \Gamma_1 = \Gamma_2 = T, \\ 0, & \Gamma_1 = A, \Gamma_2 = P. \end{cases}$$

In other words, the quark-diquark approximation of the three-quark structure of baryon octet can be performed for tensor and vector three-quark currents, independently. The preliminary consideration of a nucleon physics in the framework of this approach [17] has shown that the tensor current is more preferable from the point of view of the best description of experimental data. Therefore, we will use only the tensor current in this work.

The Ward identity does not impose any restrictions on $C_{\Gamma_1\Gamma_2}$ in the case of the baryon octet. Thus, we have only two free parameters Λ_D and C_{VT}. Their numerical values were defined by fitting experimental data and turned out to be equal to

$$\Lambda_D = 827.7 Mev, \qquad C_{VT} = \frac{3}{4} \ .$$

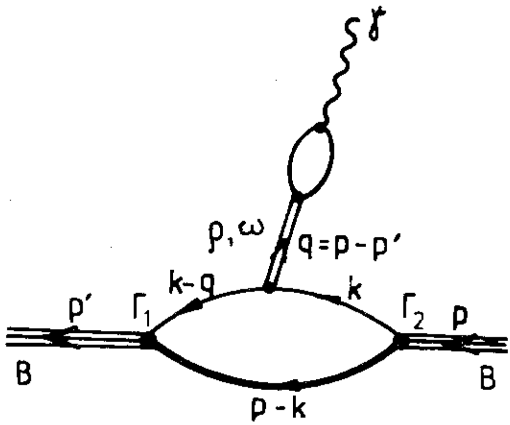
5. Basic Properties of Baryons in the QCM

Let us discuss the numerical results obtained in our approach for the basic baryon characteristics. We have calculated magnetic moments of the baryon octet and of $D \to B + \gamma$-transition, electromagnetic radii and form factors of a nucleon, the ratio of the axial and vector constants in the nonleptonic decays of a baryon octet, strong meson-nucleon constants and decay widths of a baryon decuplet.

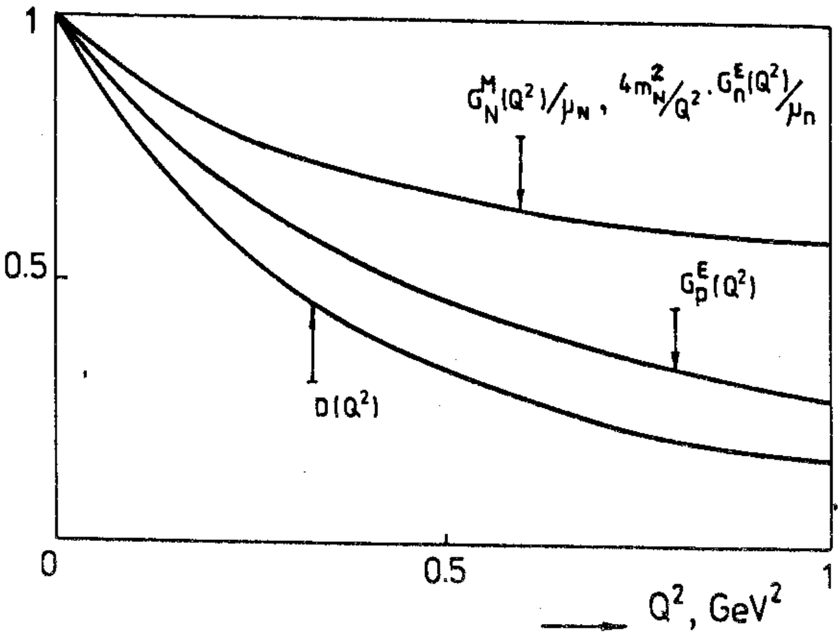
Electromagnetic baryon characteristics as magnetic moments, radii and form factors are defined by the triangle diagram, Fig.2b, and resonance one, Fig.3, taking into account intermediate vector mesons $V = \rho, \omega$.

The matrix element corresponding to the process $B \to B + \gamma$ on the baryon mass shell is written in the form

$$\begin{aligned} M(B \to B + \gamma) &= eA_\mu [F_0(m_B, q^2) tr\{\bar{B}\gamma_\mu [Q, B]\} - \\ &- \frac{i}{2m_B} F_1(m_B, q^2) tr\{\bar{B}\sigma^{\mu\nu} q_\nu [Q, B]\} - \\ &- \frac{i}{2m_B} F_2(m_B, q^2) tr\{\bar{B}\sigma^{\mu\nu} q_\nu \{Q, B\}\}] \ , \end{aligned}$$



3. Electromagnetic vertex with ρ and ω resonances.



4. Nucleon electomagnetic form factors.

where $\sigma^{\mu\nu} = \frac{i}{2}[\gamma^\mu\gamma^\nu - \gamma^\nu\gamma^\mu]$, $[Q,B] = QB - BQ$, $\{Q,B\} = QB + BQ$. Here, $F_0(m_B,q^2)$ is electic form factor of the baryon octet normalized to unity $F_0(m_B,0) = 1$. $F_1(m_B,q^2)$ and $F_2(m_B,q^2)$ are magnetic form factors. In the QCM the form factors $F_i(m_B,q^2)$, i=1,2 at the point $q^2 = 0$ are represented as

$$F_1(m_B,0) = \frac{\Phi_1(m_B)}{\Phi_0(m_B)}, \qquad F_2(m_B,0) = \frac{\Phi_2(m_B)}{\Phi_0(m_B)} ,$$

where

$$\Phi_0(m_B) = \int\limits_0^1 d\alpha \frac{\alpha(1-\alpha)a(\omega_B)}{1+\alpha((\Lambda_D/\Lambda_q)^2-1)},$$

$$\Phi_1(m_B) = \int\limits_0^1 d\alpha \frac{(1-\alpha)(2/3-\alpha)a(\omega_B)}{1+\alpha((\Lambda_D/\Lambda_q)^2-1)},$$

$$\Phi_2(m_B) = \int\limits_0^1 d\alpha \frac{(1-\alpha)a(\omega_B)}{1+\alpha((\Lambda_D/\Lambda_q)^2-1)},$$

$$\omega_B = -\frac{m_B^2\alpha(1-\alpha)}{1+\alpha((\Lambda_D/\Lambda_q)^2-1)} .$$

The magnetic moments of a baryon octet are expressed through F_1 and F_2 in a standard manner:

$$\mu_j = \begin{cases} 1+F_1(m_j,0)+\frac{1}{3}F_2(m_j,0), & j = p,\Sigma^+ \\ -\frac{2}{3}F_2(m_j,0), & j = n,\Xi^0 \\ -1-F_1(m_j,0)+\frac{1}{3}F_2(m_j,0), & j = \Sigma^-,\Xi^- \\ \frac{1}{3}F_2(m_j,0), & j = \Sigma^0 \\ -\frac{1}{3}F_2(m_j,0), & j = \Lambda^0 \end{cases} .$$

The magnetic moment of the transition $\Sigma^0 \to \Lambda^0 + \gamma$ is written as

$$\mu_{\Sigma^0\Lambda^0} = \frac{\Phi_3(m_{\Sigma^0},m_{\Lambda^0})}{2\sqrt{3\Phi_0(m_{\Sigma^0})\Phi_0(m_{\Lambda^0})}}\left(\sqrt{\frac{m_{\Sigma^0}}{m_{\Lambda^0}}} + \sqrt{\frac{m_{\Lambda^0}}{m_{\Sigma^0}}}\right) ,$$

where

$$\Phi_3(m_{\Sigma^0},m_{\Lambda^0}) = \int\limits_0^1 d\alpha \int\limits_0^1 d\beta \frac{(1-\alpha)a(\omega_{\Sigma^0\Lambda^0})}{1+\alpha((\Lambda_D/\Lambda_q)^2-1)},$$

$$\omega_{\Sigma^0\Lambda^0} = -\frac{\alpha(1-\alpha)(\beta m_{\Sigma^0}^2 + (1-\beta)m_{\Lambda^0}^2)}{1+\alpha((\Lambda_D/\Lambda_q)^2 - 1)} \quad .$$

One has to remark that the magnetic moments of a baryon octet satisfy the SU(6)-relations when the baryon masses are equal to one another:

$$\mu_p = \mu_{\Sigma^+}, \qquad \mu_n = \mu_{\Xi^0} = -\frac{2}{3}\mu_p,$$

$$\mu_{\Sigma^-} = \mu_{\Xi^-} = \mu_{\Lambda^0} = -\frac{1}{3}\mu_p,$$

$$\mu_{\Sigma^0} = \frac{1}{3}\mu_p, \qquad \mu_{\Sigma^0\Lambda^0} = \frac{1}{\sqrt{3}}\mu_p.$$

The nucleon electromagnetic form factors are parameterized as

$$G_N^E(Q^2) = F_N^1(Q^2) - \frac{Q^2}{4m_N^2}F_N^2(Q^2),$$

$$G_N^M(Q^2) = F_N^1(Q^2) + F_N^2(Q^2), \qquad Q^2 = -q^2, \qquad N = p, n \ ,$$

where

$$F_N^1(Q^2) = F_0(m_N, Q^2), \qquad F_N^2(Q^2) = F_1(m_N, Q^2) + \frac{1}{3}F_2(m_N, Q^2).$$

The experimental data are described quite accurately by the empirical dipole formula

$$G_p^E(Q^2) \simeq \frac{G_N^M(Q^2)}{\mu_N} \simeq \frac{4m_n^2}{Q^2}\frac{G_n^E(Q^2)}{\mu_n} \simeq D(Q^2),$$

where

$$D(Q^2) = \frac{1}{[1 + (Q^2/0.71Gev^2)]^2} \quad .$$

The electromagnetic radii are defined as

$$< r_p^2 >^E \; = \; -6\frac{(G_p^E(0))'}{G_p^E(0)}, \quad < r_n^2 >^E = -6(G_n^E(0))'$$

$$< r_N^2 >^M \; = \; -6\frac{(G_N^M(0))'}{G_N^M(0)}.$$

The magnetic moment μ^* of the D-B-transition is defined by non-diagonal matrix element of the operator of the electromagnetic current between the baryon octet B and decuplet D states. The matrix element of this transition is writen in the form

$$M(D \to B+\gamma) = \frac{e}{m_p}\frac{\sqrt{3}}{2}\bar{B}^k_r(p')\gamma^5(\gamma^\nu V_1(m_B,m_D)+p'_\nu V_2(m_B,m_D))$$
$$(\lambda_i Q)^{rs}\lambda_i^{nt}D^{smt}_\mu(p)F_{\mu\nu}(q)\varepsilon^{kmn},$$

where $F_{\mu\nu}(q) = q_\mu A_\nu(q) - q_\nu A_\mu(q)$, V_1 and V_2 are the vertex functions defined by a very cumbersome formula so that they are not shown here. The magnetic moment of the D-B transition is defined by the following formulae:

$$\mu^*_{DB} = \frac{1}{3}\sqrt{\frac{m_D}{m_B}}\{-V_1(m_B,m_D)(3+\frac{m_B}{m_D})+$$
$$+ 2V_2(m_B,m_D)(1-\frac{m_B}{m_D})\} .$$

It is convenient to represent the quantity μ^*_{DB} in the form

$$\mu^*_{DB} = C_{DB}\frac{2\sqrt{2}}{3}\mu_p ,$$

where μ_p is the proton magnetic moment.

The obtained results are shown in Table 2. One can see, there is only a qualitative agreement of our results with experimental data. It has to be remarked that the numerical values of magnetic moments are closed to the SU(6)-model predictions.

The dependence of the electromagnetic nucleon form factors on the square of space-like momentum $Q^2 = -q^2$ in the interval $0 \le Q^2 \le 1Gev^2$ is shown in Fig.4. One can see, there is only a qualitative agreement with the dipole formula.

The semileptonic weak decays of the baryon octet $B' \to B + e + \bar{\nu}$ are described by the triangle diagram in Fig.2b. The matrix element is written as

Table 2
Electomagnetic characteristics of baryons

Process	Observable value	Expt. [23,27]	SU_6	QCM	Other approaches [27]-[29]
$p \to p\gamma$	μ_p	2.793	3	3.660	2.79
	r_p^E, fm	0.862±0.012		0.682	0.810
	r_p^M, fm	0.858±0.056		0.560	0.810
$n \to n\gamma$	μ_n	-1.913	-2	-2.440	-1.86
	$< r_n^2 >^E$, fm^2	-0.117±0.002		-0.162	-0.130
	r_n^M, fm	0.876±0.070		0.560	0.810
$\Sigma^+ \to \Sigma^+\gamma$	μ_{Σ^+}	2.42±0.05	3	3.727	2.68
$\Sigma^0 \to \Sigma^0\gamma$	μ_{Σ^0}		1	1.242	
$\Sigma^- \to \Sigma^-\gamma$	μ_{Σ^-}	-1.157±0.025	-1	-1.243	-1.05
$\Lambda^0 \to \Lambda^0\gamma$	μ_{Λ^0}	-0.613±0.004	-1	-1.230	-0.58
$\Sigma^0 \to \Lambda^0\gamma$	$\mu_{\Sigma^0\Lambda^0}$	1.61±0.08	$\sqrt{3}$	2.217	
$\Xi^- \to \Xi^-\gamma$	μ_{Ξ^-}	-0.69±0.04	-1	-1.271	-0.47
$\Xi^0 \to \Xi^0\gamma$	μ_{Ξ^0}	-1.25±0.014	-2	-2.539	-1.40
$\Delta^+ \to p\gamma$	C_{Δ^+p}	1.25±0.2	1	1.17	
$\Delta^0 \to n\gamma$	C_{Δ^0p}	1.25±0.2	1	1.17	
$\Sigma^{*+} \to \Sigma^+\gamma$	$C_{\Sigma^{*+}\Sigma^+}$	-(1.25±0.2)	-1	-0.94	
$\Sigma^{*0} \to \Sigma^0\gamma$	$C_{\Sigma^{*0}\Sigma^0}$	0.63±0.1	0.5	0.47	
$\Sigma^{*0} \to \Lambda^0\gamma$	$C_{\Sigma^{*0}\Lambda^0}$	-(1.08±0.17)	$-\frac{\sqrt{3}}{2}$	-0.87	
$\Xi^{*0} \to \Xi^0\gamma$	$C_{\Xi^{*0}\Xi^0}$	-(1.25±0.2)	-1	-0.89	

$$M(B' \to B + e + \bar{\nu}) = \frac{G_F}{\sqrt{2}} l_\mu(q) \bar{B}'(p') \Lambda^\mu_{weak}(p, p') B(p) =$$
$$= \frac{G_F}{\sqrt{2}} l_\mu F_{weak}(m_B, m_{B'}, q^2) \{F^V_\mu + \beta[(1-\alpha)F^A_\mu + \alpha D^A_\mu]\} =$$
$$= \frac{G_F}{\sqrt{2}} F_{weak}(m_B, m_{B'}, q^2) l_\mu \bar{\psi}_{B'} [\gamma^\mu G_V - \gamma^\mu \gamma^5 G_A] \psi_B \ ,$$

where

$$F^V_\mu = tr(\bar{B}' \gamma_\mu [J, B]), \quad F^A_\mu = tr(\bar{B}' \gamma_\mu \gamma_5 [J, B]),$$

$$F^V_\mu = tr(\bar{B}' \gamma_\mu \gamma_5 \{J, B\}), \quad B = \frac{1}{\sqrt{2}} tr(\lambda \psi_B).$$

Here, $F_{weak}(m_B, m_{B'}, 0)$ is equal to

$$F_{weak}(m_B, m_{B'}, 0) = \frac{\Phi_4(m_B, m_{B'})}{2\sqrt{\Phi_0(m_B)\Phi_0(m_{B'})}} \left(\sqrt{\frac{m_B}{m_{B'}}} + \sqrt{\frac{m_{B'}}{m_B}}\right),$$

where

$$\Phi_4(m_B, m'_B) = \int\limits_0^1 d\alpha \int\limits_0^1 d\beta \frac{\alpha(1-\alpha) a(\omega_{BB'})}{1 + \alpha((\Lambda_D/\Lambda_q)^2 - 1)},$$

$$\omega_{BB'} = -\frac{\alpha(1-\alpha)(\beta m^2_{B'} + (1-\beta) m^2_B)}{1 + \alpha((\Lambda_D/\Lambda_q)^2 - 1)} \ .$$

The obtained results are shown in Table 3. It has to be remarked that the numerical values of the ratios G_A/G_V and parameters α and β coincide with the SU(6)-model predictions. Thus, there is only a qualitative agreement of our results with experimental data.

The strong meson-baryon interactions are described by the diagram in Fig.1b. The vertex functions can be written in the following way on the baryon mass shell.

1. The Meson-nucleon form factors.

a) Pseudoscalar mesons $P(\pi, \eta, \eta')$:

$$\Lambda_{PNN}(q^2) = T_P i\gamma^5 G_{PNN}(q^2) \qquad T_\pi = \vec{\tau}, \qquad T_\eta = T'_\eta = I.$$

Table 3
Weak coupling constant ratious

Process	Observable value	Expt. [23]	SU_6	QCM
$n \to pe\bar{\nu}$	G_A/G_V	-1.259±0.004	-5/3	-5/3
$\Lambda^0 \to pe\bar{\nu}$	G_A/G_V	-0.696±0.025	-1	-1
$\Sigma^- \to ne\bar{\nu}$	$\|G_A/G_V\|$	0.36±0.05	1/3	1/3
$\Xi^- \to \Lambda^0 e\bar{\nu}$	G_V/G_A	-0.25±0.05	-1/3	-1/3
$\Sigma^- \to ne\bar{\nu}$	G_A/G_V	0.01±0.10	0	0
$\Xi^- \to \Sigma^0 e\bar{\nu}$	G_A/G_V		-5/3	-5/3
$\Xi^0 \to \Sigma^+ e\bar{\nu}$	G_A/G_V		-5/3	-5/3
	β	-1.259±0.004	-5/3	-5/3
	α	0.66±0.07	3/5	3/5

Table 4
Strong meson-nucleon constants

Vertex	$G^2_{MNN}(0)/(4\pi)$	
	QCM	Other approaches[30]
πNN	12.32	14.08; 14.28 ± 0.018
ηNN	11.21	3.67; 5.0
$\eta' NN$	7.99	3.77; 4.23
$a_0 NN$	0.98	1.62; 1.16
εNN	13.93	4.56; 8.85
ρNN	0.42 (F/G=5.1)	0.41; 0.55 ± 0.06 (F/G=6.1) (F/G=6.1 ±0.6)
ωNN	3.78 (F/G=0.22)	10.6; 5.7 ± 2.0 (F/G=0) (F/G=0)

b) Scalar mesons $S(a_0, \varepsilon)$:

$$\Lambda_{SNN}(q^2) = T_S G_{SNN}(q^2) \qquad T_{a_0} = \vec{\tau}, \qquad T_\varepsilon = I.$$

c) Vector mesons $V(\rho, \omega)$:

$$\Lambda_{VNN}(q^2) = T_V\{\gamma^\mu G_{VNN}(q^2) - \frac{i}{2m_N}\sigma^{\mu\nu} q_\nu F_{VNN}(q^2)\} \ .$$

$$T_\rho = \vec{\tau}, \qquad T_\omega = I \ .$$

2. Decay $D \to B + \pi$

$$\Lambda^\mu_{\pi BD}(q^2) = q^\mu \frac{G_{\pi BD}(q^2)}{m_\pi} \ .$$

The numerical values of the $G^2_{MNN}(0)/4\pi$ and $F_{VNN}(0)/G_{VNN}(0)$ are shown in Table 4. It also shows the results of phenomenological approches [27,30]. One can see, the results for the πNN-, ρNN-, ωNN- and $a NN$-form factors obtained in the QCM coincide with the phenomenological ones. The vector meson-nucleon constants are connected by SU(3)-relations and are immediately defined by the coupling constant of vector mesons with quarks g_V and nucleon magnetic moments:

$$\frac{G^2_{\rho NN}(0)}{4\pi} = \frac{g_v^2}{8\pi}, \quad G_{\omega NN}(0)/G_{\rho NN}(0) = 3,$$

$$F_{\rho NN}(0)/G_{\rho NN}(0) = \mu_p - \mu_n - 1, \quad F_{\omega NN}(0)/G_{\omega NN}(0) = \mu_p + \mu_n - 1.$$

It is interesting to consider the case when all hadron masses are equal to zero. We have

$$L_{eff} = G_{\rho NN}\bar{N}\gamma^\mu\vec{\tau}N\vec{\rho}_\mu + G_{\rho qq}\bar{q}\gamma^\mu\vec{\tau}q\vec{\rho}_\mu + G_{\rho\pi\pi}\rho^i\pi^j\partial^\mu\pi^k\varepsilon^{ijk} \ ,$$

where

$$G_{\rho NN} = G_{\rho qq} = G_{\rho\pi\pi} = 1/g_{\rho\gamma} = \pi.$$

(Here, we accept a standard notation for the effective coupling constants instead of ones used in our model: $G_{\rho qq} = g_\rho/\sqrt{2}$ and $G_{\rho\pi\pi} = g_{\rho\pi\pi}/2$.)

Table 5
Decay $D \to B + \pi$ widths

Processes	Expt.[23], $\Gamma(Mev)$	QCM, $\Gamma(Mev)$
$\Delta^{++} \to p + \pi^+$	111.5±0.67	111.6
$\Delta^{+} \to p + \pi^0$	76.1±0.46	74.4
$\Delta^{+} \to n + \pi^+$	36.8±0.22	37.2
$\Delta^{0} \to p + \pi^-$	38.3±0.23	37.2
$\Delta^{0} \to n + \pi^0$	76.7±0.46	74.4
$\Delta^{-} \to n + \pi^-$	116±0.69	111.6
$\Xi^{*0} \to \Xi^0 + \pi^0$	3.26±0.2	4.66
$\Xi^{*0} \to \Xi^- + \pi^+$	5.54±0.25	8.05
$\Xi^{*-} \to \Xi^- + \pi^0$	3.09±0.2	4.32
$\Xi^{*-} \to \Xi^0 + \pi^-$	6.56±0.4	9.94
$\Sigma^{*+} \to \Lambda^0 + \pi^+$	30.8±0.61	35.67
$\Sigma^{*+} \to \Sigma^+ + \pi^0$	2.3±0.046	2.69
$\Sigma^{*+} \to \Sigma^0 + \pi^+$	1.9±0.04	2.43
$\Sigma^{*0} \to \Sigma^- + \pi^+$	1.86±0.037	2.14
$\Sigma^{*0} \to \Sigma^+ + \pi^-$	2.3±0.046	2.77
$\Sigma^{*0} \to \Lambda^0 + \pi^0$	30.8±0.61	36.15
$\Sigma^{*-} \to \Lambda^0 + \pi^-$	35.2±0.71	37.61
$\Sigma^{*-} \to \Sigma^0 + \pi^-$	2.4±0.048	2.75

This result is in complete coincidence with the $\rho-$ universality hypothesis and the prediction of an effective gauge field theory [31].

The decay $D \to B + \pi$ width is calculated according to the formula:

$$\Gamma(D \to B + \pi) = m_\pi \frac{G^2_{\pi BD}(m_\pi^2)}{12\pi} (\frac{m_B}{m_D})(\frac{p^*}{m_\pi}^3)(1 + \frac{E_B}{m_B}),$$

where

$$p^* = \frac{\sqrt{(m_D^2 - (m_B + m_\pi)^2)(m_D^2 - (m_B - m_\pi)^2)}}{2m_D},$$

$$E_B = \frac{m_D^2 + m_B^2 - m_\pi^2}{2m_D}$$

are the baryon momentum energy in the c.m. system. The results are shown in Table 5. One can see, our results are in a quite good agreement with experimental data [23].

Thus, the quark-diquark approximation of the three-quark structure of a baryon, in which a diquark is considered to be a hard core, correctly reproduces the baryon inner structure at low energies and allows one to describe most of the baryon properties from a unique point of view.

In future, we plan to use this picture for the description of more interesting processes of baryon physics as the baryon nonleptonic decays, photoproduction of π and η mesons on nucleons, NN-scattering phase shifts, etc.

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Appendix

To demonstrate the calculational technique in the QCM, we calculate the meson mass operator and the baryon vertex.

The meson mass operator is written as

$$\Pi^{\Gamma_1\Gamma_2}(p)=\int\frac{d^4k}{4\pi^2 i}\int d\sigma_v tr[\Gamma_1\frac{1}{v\Lambda_q-\hat{k}-\hat{p}}\Gamma_2\frac{1}{v\Lambda_q-\hat{k}}]\,.$$

By using the Feynman α -parametrization one can obtain

$$\Pi^{\Gamma_1\Gamma_2}(p)=\int\limits_0^1 d\alpha\int\limits_0^\infty du\int d\sigma_v\frac{R(u,\alpha,p)}{[v^2+u-\alpha(1-\alpha)p^2/\Lambda_q^2]^2}\,.$$

Here

$$R(u,\alpha,p)=\frac{1}{4}tr[v^2\Lambda^2\Gamma_1\Gamma_2-\alpha(1-\alpha)\Gamma_1\hat{p}\Gamma_2\hat{p}-\frac{u}{4}\Lambda_q^2\Gamma_1\gamma^\alpha\Gamma^2\gamma_\alpha]+$$
$$+v\Lambda_q(1-\alpha)\Gamma_1\hat{p}\Gamma_2-\alpha\Gamma_1\Gamma_2\hat{p}]\,.$$

Recalling the definitions of the confinement function $G(z)$, we have

$$\int_0^\infty duu \int d\sigma_v \frac{v^2}{[v^2+u+\Delta(\alpha)]^2} = -\int_0^\infty du(u+\Delta(\alpha))b(u+\Delta(\alpha)),$$

$$\int_0^\infty duu \int d\sigma_v \frac{1}{[v^2+u+\Delta(\alpha)]^2} = \int_0^\infty dub(u+\Delta(\alpha)),$$

$$\int_0^\infty duu^2 \int d\sigma_v \frac{1}{[v^2+u+\Delta(\alpha)]^2} = 2\int_0^\infty dub(u+\Delta(\alpha)),$$

where

$$\Delta(\alpha) = -\frac{\alpha(1-\alpha)p^2}{\Lambda_q^2} \quad .$$

Using these formulas one can obtain

$$\Pi^{\Gamma_1\Gamma_2}(p) = -\frac{\Lambda_q^2}{4}tr[B_1(s)(\Gamma_1\Gamma_2+\frac{1}{2}\Gamma_1\gamma^\alpha\Gamma_2\gamma_\alpha)+$$
$$+\frac{1}{6}B_0(s)(\Gamma_1\frac{\hat{p}}{\Lambda_q}\Gamma_2\frac{\hat{p}}{\Lambda_q}+2s\Gamma_1\gamma^\alpha\Gamma_2\gamma_\alpha)+\frac{1}{2}A_0(s)(\Gamma_1\Gamma_2-\Gamma_2\Gamma_1)\frac{\hat{p}}{\Lambda_q}],$$

$$s = \frac{p^2}{4\Lambda_q^2}$$

Here

$$B_1(s) = \int_0^\infty duub(u) - s^2\int_0^1 duub(-us)\sqrt{1-u},$$

$$B_0(s) = \int_0^\infty dub(u) + s\int_0^1 dub(-us)\sqrt{1-u}(1+\frac{u}{2}),$$

$$A_0(s) = \int_0^\infty dua(u) + s\int_0^1 dua(-us)\sqrt{1-u}.$$

The baryon vertex is written as

$$\Lambda_\Gamma(p,p') = d^{\Gamma_1\Gamma_2}\Gamma_1 T_\Gamma(p,p')\Gamma_2,$$

$$T_\Gamma(p,p') = \int \frac{d^4k}{\pi^2 i}\int d\sigma_v \frac{1}{v\Lambda_q - \hat{k}+\hat{q}}\Gamma\frac{1}{v\Lambda_q-\hat{k}}\frac{1}{v^2\Lambda_D^2-(p-k)^2}.$$

Making use of the Feynman α-parametrization, one can obtain

$$T_\Gamma(p,p') = \int\limits_0^1 d\mu_\alpha \int\limits_0^\infty duu \int d\sigma_v \frac{F(u,\alpha,p,p')}{[v^2+u+\Delta(\alpha)]^3},$$

$$d\mu_\alpha = 2d^3\alpha\delta(1-\alpha_1-\alpha_2-\alpha_3).$$

where

$$F(\alpha,u,p,p') = -\frac{1}{4}\gamma^\alpha\Gamma\gamma_\alpha u +$$
$$+\frac{1}{1+\alpha((\frac{\Lambda_D}{\Lambda_q})^2-1)}(\Gamma v^2 + v\frac{\hat{r}_1\Gamma+\Gamma\hat{r}_2}{\Lambda_q}+\frac{\hat{r}_1\Gamma\hat{r}_2}{\Lambda_q^2}),$$

$$\hat{r}_1 = \hat{p}'(1-\alpha_2)-\hat{p}\alpha_3, \quad \hat{r}_2 = \hat{p}(1-\alpha_3)-\hat{p}'\alpha_2,$$

$$\Delta(\alpha) = -\frac{1}{1+\alpha(((\Lambda_D/\Lambda_q))^2-1)}(p'^2\alpha_1\alpha_2+p^2\alpha_1\alpha_3+q^2\alpha_2\alpha_3).$$

Recalling the definition of the confinement function we have

$$\int\limits_0^\infty duu\int d\sigma_v\frac{v^2}{[v^2+u+\Delta(\alpha)]^3} = -\frac{\Delta(\alpha)}{2}b(\Delta(\alpha)),$$

$$\int\limits_0^\infty duu^2\int d\sigma_v\frac{1}{[v^2+u+\Delta(\alpha)]^3} = \int\limits_0^\infty dub(u+\Delta(\alpha)),$$

$$\int\limits_0^\infty duu\int d\sigma_v\frac{1}{[v^2+u+\Delta(\alpha)]^3} = \frac{1}{2}b(\Delta(\alpha))$$

$$\int_0^\infty duu \int d\sigma_v \frac{v}{[v^2+u+\Delta(\alpha)]^3} = \frac{1}{2}a(\Delta(\alpha))$$

So, the final results is written as

$$\begin{aligned} T_\Gamma(p,p') &= -\frac{1}{4}[\int_0^\infty dub(u) - \int d\mu_\alpha \int_0^{\Delta(\alpha)} dub(u)]\gamma^\alpha\Gamma\gamma_\alpha + \\ &+ \frac{1}{2}\int \frac{d\mu_\alpha}{1+\alpha((\Lambda_D/\Lambda_q)^2-1)}[a(\Delta(\alpha))\frac{\hat{r}_1\Gamma+\Gamma\hat{r}_2}{\Lambda_q} + \\ &+ b(\Delta(\alpha))(-\Delta(\alpha)\Gamma + \frac{\hat{r}_1\Gamma\hat{r}_2}{\Lambda_q^2})]. \end{aligned}$$

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