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E2-88-37

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CONFINEMENT AND QUARK STRUCTURE OF LIGHT HADRONS

Submitted to "International Journal
of Modern Physics"

1988

1. INTRODUCTION

The world of the light quarks and hadrons can schematically be represented as follows. At short distances (up to 0.2-0.3 fm) there are only free quarks and gluons. At large distances (over 1 fm) there are only hadrons. The quarks and gluons are governed by quantum chromodynamics (QCD). At short distances the running quark-gluon coupling constant is small enough for perturbative theory to apply a property of QCD known as asymptotic freedom. At large distances the point-like hadrons are described by the standard quantum field equations. At intermediate distances (0.2-1 fm) colour confinement and hadronization take place. We are still far from the theoretical explanation of these phenomena[1]. From the physical point of view this is a low energy region of hadron physics where physical processes with the liberated energy 1-2 GeV proceed. The most nonperturbative QCD methods slip through this region and consider only the connection of the hadron amplitudes with the quark-gluon parameters. The confinement problem is passed by anyhow.

For example, QCD sum rules[2] connect matrix elements of hadron currents as dispersion integrals with the corresponding quark-gluon diagrams. The nonperturbative effects are taken into account by introducing quark and gluon condensates.

The bosonization of QCD[3] aims to obtain phenomenological meson Lagrangians by using the gauge and chiral groups escaping the problems such as the dynamical hadronization and confinement.

In the papers[4] we have developed the so-called virton-quark model (VQM) based on a specific notion of the quark and hadron behaviour in the confinement region. The quarks in a low-energy region were assumed to be like quasiparticles having quantum quark numbers but described by the virton field existent in the virtual state only. This assumption is provided by the requirement that the quark propagator is chosen to be an entire analytical function on a complex plane of the momentum p^2. Thus, from the beginning the requirement of confinement is satisfied in the VQM. Further, it is suggested that the coupling of hadrons with quark-virtons can be given by the corresponding interaction Lagrangians. The fact that hadrons consist of quarks is taken into account by the quantum-field compositeness condition defining the hadron-quark coupling constants. The knowledge of the transition dynamics of a hadron into quarks and vice versa allows us to describe all possible physical processes at the low energy. However, at this wording, the VQM has no connection with QCD.

Now we will present the quark confinement model (QCM) based on the QCD ideas. In particular, it will be shown under what assumptions the VQM can be obtained from the QCM.

The present model is based on the physical picture which is believed to be valid at the intermediate distances. A complex and nontrivial gluon vaccuum of QCD is created. There are many papers/5/ devoted to this problem but we are still far from solving it. This gluon vacuum is the kind of substance in which no colour objects exist as free particles. In other words, the gluon vacuum of QCD provides the colour confinement.

The hadrons as colourless states appear to be the collective variables corresponding to the colourless quark currents. This idea was developed in many papers/6/.

The mathematical realization of this picture in terms of the hadron scattering matrix requires some hypothesis about, first, the confinement as a way of averaging over the vacuum gluon fields for all quark diagrams and second, the hadronization as a transition to collective variables in the QCD Lagrangian. The assumptions of this kind form the contents of the present quark confinement model (QCM).

The VQM can be obtained from QCM as one of the possibilities of averaging over the gluon vacuum.

The QCM is applied to description of the low energy physics. The main characteristics of light mesons are calculated and good agreement with experimental data is obtained.

2. QCD VACUUM HYPOTHESIS AND HADRONIZATION

Let us adduce arguments on which our model is based. We proceed from the QCD Lagrangian

$$\mathcal{L}_{QCD} = -\frac{1}{8g^2}\,\mathrm{tr}\, G_{\mu\nu}^2 + \sum_{f=1}^{3} \bar{q}_f(x)\left(i\hat{\partial} - m_f - \hat{B}(x)\right) q_f(x). \tag{2.1}$$

Here $\hat{B} = B_\mu^a \gamma_\mu t^a$, $G_{\mu\nu} = \partial_\mu B_\nu - \partial_\nu B_\mu + [B_\mu, B_\nu]$, B_μ^a is the gluon field and t^a are the matrices of a colour group SU (3), q_f (f = 1,2,3) are the quark fields with mass m_f, g is the QCD coupling constant.

Let us consider the QCD functional

$$Z = \iiint \delta q\, \delta \bar{q}\, \delta B\, \mathcal{D}[B] \exp\left\{ i \int dx\, \mathcal{L}_{QCD}(x) \right\}, \tag{2.2}$$

where $\mathcal{D}[B]$ allows for the choice of a gauge. The functional (2.2) is hoped to answer such hard questions as the colour confinement and the hadronization.

The accepted point of view is that these phenomena are provided by the complex structure of gluon vacuum. One can imagine, the QCD vacuum is the substance in which only colourless objects can exist, but the colour ones cannot exist as free particles.

We suppose that the QCD vacuum is realized on the quantum gluon fields with indefinite metrics. We will mark these fields $B_{vac}(x)$. The gluon vacuum is supposed to be degenerated, e.g., there is a set of gluon vacuum fields characterized by the independent parameters $\{c_{vac}\}$, which are the degrees of freedom of vacuum. Thus, the gluon field can be represented as

$$B(x) = B_{vac}(x) + B_{fl}(x),$$

where B_{fl} are the quantum fluctuations around the QCD vacuum. These fields depend on B_{vac} and are defined by the independent parameters $\{c_{fl}\}$. The functional differential can be written as

$$\delta B = \prod_x dB(x) = \prod dc_{vac} \prod dc_{fl}\, \mathcal{D}(B \,|\, c_{vac}, c_{fl}).$$

were assumed to be like quasiparticles having quantum quark numbers but described by the virton field existent in the virtual state only. This assumption is provided by the requirement that the quark propagator is chosen to be an entire analytical function on a complex plane of the momentum p^2 . Thus, from the beginning the requirement of confinement is satisfied in the VQM. Further, it is suggested that the coupling of hadrons with quark-virtons can be given by the corresponding interaction Lagrangians. The fact that hadrons consist of quarks is taken into account by the quantum-field compositeness condition defining the hadron-quark coupling constants. The knowledge of the transition dynamics of a hadron into quarks and vice versa allows us to describe all possible physical processes at the low energy. However, at this wording, the VQM has no connection with QCD.

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$$\delta B = \prod_x dB(x) = \prod dc_{vac} \prod dc_{fl}\, \mathcal{D}(B \mid c_{vac}, c_{fl}),$$

where $\mathcal{D}(B|c_{vac}, c_{f\ell})$ is determinant of transformation into the variables $\{c_{vac}\}$ and $\{c_{f\ell}\}$.

Substituting these formulae into (2.2) and integrating over the $\{c_{f\ell}\}$, we have

$$Z = \int d\sigma_{vac} \iint \delta q\, \delta\bar{q}\, \exp\Big\{ i \int dx \sum_f \bar{q}_f (i\hat{\partial} - m_f - \hat{B}_{vac}) q_f + \sum_n L_n \Big\}, \tag{2.3}$$

$$L_n = g^n \int dx_1 \ldots \int dx_n \prod_{j=1}^{n} J^{a_j}_{\mu_j}(x_j)\, G^{a_1\ldots a_n}_{\mu_1\ldots\mu_n}(x_1,\ldots,x_n|B_{vac}).$$

Here

$$J^a_\mu(x) = \sum_f \bar{q}_f(x)\gamma_\mu t^a q_f(x)$$

is the colour quark current and

$$d\sigma_{vac} = \prod d c_{vac}\, F[B_{vac}] \tag{2.4}$$

is the indefinite measure of integration over the gluon vacuum fields. The shape of this measure must be determined by solving the QCD vacuum problem.

The functions $G^{a_1\ldots a_n}_{\mu_1\ldots\mu_n}(x_1,\ldots,x_n|B_{vac})$ are the full n-point connected Green functions of gluons in the vacuum field B_{vac}.

Suppose that at the large distances these Green functions can be changed by

$$G^{a_1\ldots a_n}_{\mu_1\ldots\mu_n}(x_1,\ldots,x_n|B_{vac}) \rightarrow \tag{2.5}$$

$$\rightarrow G^{a_1\ldots a_n}_{\mu_1\ldots\mu_n}(x_1,\ldots,x_n) = \int d\sigma_{vac}\, G^{a_1\ldots a_n}_{\mu_1\ldots\mu_n}(x_1,\ldots,x_n|B_{vac}).$$

where $G^{\{a\}}_{\{\mu\}}(x)$ is the n-point connected Green function of gluon fields in the strong coupling regime. This approximation denotes that the connections of the quark fields with the different gluon Green functions by the gluon vacuum can be neglected in the functional (2.3).

The next step is the hadronization. By the hadronization we shall imply the transition to the collective variables in the representation

(2.3). The idea is to pass from the colour quark currents $J^a_\mu(x)$ in each term L_n to the product of colourless quark objects. Other words, the set of colourless hadron states which are made of n quarks must be produced by the term L_n.

Let us demonstrate the appearance of the two quark states (mesons). For this purpose in (2.3) only the term L_2 is taken into account. We have

$$L_2 = \iint dx_1 dx_2\, J^{a_1}_{\mu_1}(x_1)\, G^{a_1 a_2}_{\mu_1\mu_2}(x_1-x_2)\, J^{a_2}_{\mu_2}(x_2),$$

where the quark current $J^a_\mu(x)$ is an octet of colour group $SU_c(3)$ and a singlet of flavour SU(3).

The gluon Green function can be written in the form

$$G^{a_1a_2}_{\mu_1\mu_2}(x_1-x_2) = \delta_{a_1a_2} g_{\mu_1\mu_2} G(x_1-x_2) = -\delta_{a_1a_2} g_{\mu_1\mu_2} \int \frac{d^4p}{(2\pi)^4 i} \tilde{G}(p^2) e^{-ip(x_1-x_2)}.$$

There are two points of view concerning the behaviour of $\tilde{G}(p^2)$ at small p^2: the one[7] that $\tilde{G}(p^2) \sim 1/p^4$ near $p^2 = 0$ to provide the linear increased potential between quarks at large distances and the other[8] that $\tilde{G}(p^2)$ is the analytical function at the point $p^2 = 0$ and may be the entire function of the plane p^2, which provides the complete gluon confinement. We shall hold the second view point.

Note that if $G(x_1-x_2) = -iG_0\,\delta(x_1-x_2)$, then L_2 is the local four-quark Lagrangian interaction like current×current. There is a line of investigations (see, ref.[9] others therein) in which the chiral phenomenological Lagrangians are obtained from the four-quark interaction the idea was advanced in[6,10]. We will follow these ideas.

Let us transform L_2. By using the Firtz transformations L_2 can be written in the form

$$L_2 = H_2 + D_2,$$

$$H_2 = -\sum_J \sum_{f=0}^{8} a_{Jf} \iint dx_1 dx_2\, J_{Jf}(x_1,x_2)\, G(x_1-x_2)\, J_{Jf}(x_2,x_1), \tag{2.6}$$

$$D_2 = -\sum_J \sum_{f=0}^{8} b_{Jf} \iint dx_1 dx_2\, \bar{B}^{af}(x_1,x_2)\, G(x_1-x_2)\, B^{af}_{Jf}(x_2,x_1).$$

The term H_2 consists of the sum of the products of colourless quark currents

$$J_{J\ell}(x_1,x_2) = \bar{q}(x_1)\, O_J \lambda_\ell\, q(x_2)\,, \qquad (J = S, P, V, A)\,,$$

where λ_ℓ is the flavour SU_3 -matrix. The term D_2 contains the products of the diquark colour currents

$$B^{ab}_{J\ell}(x_1,x_2) = q^a(x_1)\, C O_J \lambda_\ell\, q^b(x_2)\,.$$

It is to be remarked that the coefficients $b_{J\ell}$ and $a_{J\ell}$ in (2.6) can be chosen by the different way. It will be shown that sign of these coefficients must provide the interpretation of the currents $J_{J\ell}$ as hadron states and quantity $a_{J\ell}$ must define their masses. In this paper we will consider only H_2.

Let us turn to new variables $x = \frac{1}{2}(x_1 + x_2)\,,\ \xi = \frac{1}{2}(x_1 - x_2)$. We have

$$H_2 = -\sum_{\Gamma = \{J,\ell\}} a_\Gamma \int dx \int d\xi\, J_\Gamma(x,\xi)\, G(\xi)\, J_\Gamma(x,-\xi)$$

$$J_\Gamma(x,\xi) = \bar{q}(x+\xi)\, O_J \lambda_\ell\, q(x-\xi)\,.$$

The gluon propagator $\tilde{G}(p^2)$ is by our hypothesis analytical function at the point $p^2 = 0$. Thus, the function $G(\xi)$ can be written in the form

$$G(\xi) = -i\int \frac{d^4p}{(2\pi)^4}\, \tilde{G}(p^2)\, e^{-ip\xi} = -i\int \frac{d^4p}{(2\pi)^4} \sum_{k=0}^{\infty} G_k\, (p^2)^k\, e^{-ip\xi} =$$

$$= -i \sum_{k=0}^{\infty} G_k\, (\Box_\xi)^k\, \delta(\xi)\,.$$

Substituting this expression in H_2 and performing the integration by parts, we obtain

$$H_2 = +i\sum_\Gamma a_\Gamma \sum_K G_K \sum_{2n+\ell = K} z_{n\ell} \int dx\, J^{(\Gamma,n,\ell)}_{\mu_1\ldots\mu_\ell}(x)\, J^{(\Gamma,n,\ell)}_{\mu_1\ldots\mu_\ell}(x) =$$

$$= +i\sum_Q \int dx\, \big(g_Q J_Q(x)\big)^2\,, \tag{2.7}$$

where $z_{n\ell}$ are the numerical coefficients and

$$J_Q(x) = J^{(\Gamma,n,\ell)}_{\mu_1\ldots\mu_\ell}(x) = \Box^n_\xi\, T^\ell_{\mu_1\ldots\mu_\ell}(\partial_\xi)\, \bar{q}(x+\xi)\,\Gamma\, q(x-\xi)\Big|_{\xi=0} =$$

$$= \bar{q}(x)\,\Gamma\, T^\ell_{\mu_1\ldots\mu_\ell}(\overleftrightarrow{\partial}_x)\,(\overleftrightarrow{\partial}_x \overleftrightarrow{\partial}_x)^n\, q(x)\,. \tag{2.8}$$

Here $T^\ell_{\mu_1\ldots\mu_\ell}(x)$ is the symmetric tensor with null trace

$$T^\ell_{\mu_1\ldots\mu_\ell}(a) = a_{\mu_1} a_{\mu_2} \ldots a_{\mu_\ell} - \delta_{\mu_1\mu_2} a_{\mu_3}\ldots a_{\mu_\ell} - \ldots\,,$$

This tensor describes orbital moment ℓ.

The current $J^{(\Gamma,n,\ell)}_{\mu_1\ldots\mu_\ell}(x)$ is the two-quark state with quantum numbers defined by indices $\Gamma = \{J,\ell\}$ and ℓ but not n . This means that n corresponds to the radial excitation of the two-quark state described by the current $J^{(\Gamma,n,\ell)}_{\mu_1\ldots\mu_\ell}(x)$. Orthogonal polynomials must exist to characterize these excitations. It is conceivable that these polynomials depend only on the two-quark state numbers Γ and ℓ . In this paper we do not discuss these questions in detail. We would like to emphasize that all the two-quark hadron states are contained in the representation (2.7).

Let us use the representation

$$e^{H_2} = \int \prod_Q \delta M_Q\, \exp\Big\{-\frac{1}{2}\int dx\, M_Q^2(x) + i\sum_Q g_Q \int dx\, M_Q(x) J_Q(x)\Big\}.$$

The field variables $M_Q(x) = M^{(\Gamma,n,\ell)}_{\mu_1\ldots\mu_\ell}(x)$ correspond to the meson field with arbitrary excitations. Inserting this representation into (2.3) and taking into account the term H_2 only, we have

$$Z = \int \prod_Q \delta M_Q \exp\{-\frac{i}{2}\int dx\, M_Q^2(x)\} \times$$

$$\times \iint \delta q\, \delta \bar q \int d\sigma_{vac} \exp\{ i\int dx\, \bar q_f(x)[i\hat\partial - m_f - \hat B_{vac} + M(x)]\, q_f(x)\},$$

$$M(x) = \sum_Q g_Q \Gamma_Q M_Q(x).$$

After integrating over the quark fields

$$Z = \int \prod_Q \delta M_Q \int d\sigma_{vac} \exp\{-\frac{i}{2}\int dx\, M_Q^2(x) -$$

$$-\sum_N \frac{i^N}{N}\int\cdots\int dx_1\dots dx_N\, tr[M(x_1)S(x_1,x_2|B_{vac})\dots M(x_N)S(x_N,x_1|B_{vac})]\}. \quad (2.9)$$

Here

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

Our next assumption consists in that the expression (2.9) can be written in the form

$$Z = \int \prod_Q \delta M_Q \exp\{-\frac{i}{2}\int dx\, M_Q^2(x) - \quad (2.10)$$

$$-\sum_N \frac{i^N}{N}\int dx_1\dots\int dx_N \int d\sigma_{vac}\, tr[M(x_1)S(x_1,x_2|B_{vac})\dots M(x_N,x_1|B_{vac})\}.$$

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

The measure $d\sigma_{vac}$ in (2.10), i.e. the averaging over the QCD gluon vacuum, must provide the quark confinement. The assumptions about the measure properties will be done in Sec. 4.

Let us give off the terms diagonalized over meson variables from the expression in the sum in (2.10):

$$-\frac{i}{2}\iint dx_1 dx_2 \sum_Q g_Q^2\, M_Q(x_1)\, \Pi_Q(x_1-x_2)\, M_Q(x_2),$$

$$\Pi_Q(x_1-x_2) = i\int d\sigma_{vac}\, tr\{\Gamma_Q S(x_1,x_2|B_{vac})\Gamma_Q S(x_2,x_1|B_{vac})\}$$

Combining this term with the quadratic one in (2.10), we have

$$-\frac{i}{2}\iint dx_1 dx_2 \sum_Q M_Q(x_1)[\delta(x_1-x_2) + g_Q^2\, \Pi_Q(x_1-x_2)]\, M_Q(x_2) =$$

$$= -\frac{i}{2}\sum_Q \int dp\, \tilde M_Q(p)[1 + g_Q^2\, \tilde\Pi_Q(p^2)]\, \tilde M_Q(p).$$

The equation

$$1 + g_Q^2\, \tilde\Pi_Q(m_Q^2) = 0 \quad (2.11)$$

defines the meson mass.

Let us pass to the following normalization in (2.10)

$$M_Q(x) \to M_Q(x)\Big/\sqrt{-g_Q^2\, \tilde\Pi'(m_Q^2)}\,.$$

The quadratic form over meson fields with the inclusion of (2.11) is written as

$$\frac{i}{2}\sum_Q \int dp\, \tilde M_Q(p)\, \frac{[\tilde\Pi_Q(m_Q^2) - \tilde\Pi_Q(p^2)]}{[-\tilde\Pi_Q'(m_Q^2)]}\, \tilde M_Q(p) =$$

$$= \frac{i}{2}\sum_Q \{\int dx\, M_Q(x)(\Box - m_Q^2) M_Q(x) - \iint dx_1 dx_2\, M_Q(x_1)\, G_Q^2\, \tilde\Pi_Q^{reg}(x_1-x_2)\, M_Q(x_2)\},$$

where

$$1 + G_Q^2\, \tilde\Pi_Q'(m_Q^2) = 0,$$

$$\tilde\Pi_Q^{reg}(p^2) = \tilde\Pi_Q(p^2) - \tilde\Pi_Q(m_Q^2) - \tilde\Pi_Q'(m_Q^2)(p^2 - m_Q^2).$$

Taking into account the replacements and introducing sources of meson fields, we write the functional (2.10) in the form

$$Z[J] = \int \prod_Q \delta M_Q \exp\Big\{\frac{i}{2}\int dx \sum_Q M_Q(x)(\Box - m_Q^2) M_Q(x) -$$

$$-\sum_N{}' \frac{i^{-N}}{N}\int dx_1 \dots \int dx_N \int d\sigma_{vac}\, \mathrm{tr}\big[M(x_1)S(x_1,x_2|B_{vac}) \dots M(x_N)S(x_N,x_1|B_{vac})\big] +$$

$$+ i\sum_Q \int dx\, M_Q(x) J_Q(x)\Big\}, \qquad (2.12)$$

where

$$M(x) = \sum_Q G_Q \Gamma_Q M_Q(x).$$

The prime in a sum (2.12) implies to use Π_Q^{reg} in the quadratic forms over M_Q.

A few remarks are to be done. First, the mass equation (2.11) and the transformation from (2.10) to (2.12) require definite relations between the signs of the coefficients a_r, G_n, $z_{n\ell}$ in (2.7) to provide the right normalization of physical quantities. This fact imposes restrictions on expansion (2.6).

Second, the representation (2.12) does not contain a_r, G_n, $z_{n\ell}$ in an explicit form which define the behaviour of the gluon propagator near the point $p^2 = 0$. Thus, the coupling constant is defined by the hadron mass and confinement shape.

Third, (2.11) connects the meson mass, the expansion coefficients of the confined gluon Green function at point $p^2 = 0$ and the universal confinement function arising after averaging of mass operator over vacuum gluon fields.

The generating function $Z[J]$ in (2.12) defining the meson-meson interactions 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 (2.12). Let us show that the functional $Z[J]$ in (2.12) can be written in the form

$$Z[J] = \int \prod_Q \delta M_Q \exp\Big\{\frac{i}{2}\sum_Q \int dx\, M_Q(x)[\Box - m_Q^2 + \delta m_Q^2] M_Q(x)\Big\} \times$$

$$\times \int d\sigma_{vac} \iint \delta q\, \delta\bar{q} \exp\Big\{ i\sum_f \int dx\, \bar{q}_f(x)(i\hat{\partial} - m_f - \hat{B}_{vac}) q_f(x) +$$

$$+ i\sum_{Qf} g_Q \int dx\, M_Q(x)\bar{q}_f(x)\Gamma_Q q_f(x) + i\sum_Q \int dx\, M_Q(x) J_Q(x)\Big\} \qquad (2.13)$$

under the condition that the wave-function renormalization constant of meson M_Q is equal to zero.

Indeed, let us integrate over the quark fields in (2.13) using the same assumptions about measure $d\sigma_{vac}$ and give off the quadratic term diagonalized over

$$\frac{i}{2}\int dp\, \tilde{M}_Q(p)[p^2 - m_Q^2 + \delta m_Q^2 - g_Q^2 \tilde{\Pi}_Q(p^2)] \tilde{M}_Q(p) =$$

$$= \frac{i}{2}\int dp\, \tilde{M}_Q(p)(p^2 - m_Q^2) Z_Q^{-1} \tilde{M}_Q(p) -$$

$$- \frac{i}{2}\int dp\, \tilde{M}_Q(p) g_Q^2 \tilde{\Pi}_Q^{reg}(p^2) \tilde{M}_Q(p),$$

where

$$\delta m_Q^2 = \tilde{\Pi}_Q(m_Q^2) g_Q^2$$

$$Z_Q^{-1} = 1 - g_Q^2 \tilde{\Pi}_Q'(m_Q^2).$$

Performing a replacement

$$M_Q \to Z_Q^{1/2} M_Q, \quad g_Q \to Z_Q^{-1/2} G_Q$$

one can obtain that the representation (2.13) coincides with (2.12) under the condition

$$G_Q^2 = 1 \Big/ \big(-\tilde{\Pi}_Q'(m_Q^2)\big).$$

This condition leads to that the wave-function renormalization constant of meson M_Q is equal to zero

$$Z_Q = \frac{1}{1-g_Q^2 \tilde{\Pi}'(m_Q^2)} = 1+ G_Q^2 \tilde{\Pi}_Q'(m_Q^2) = 0. \tag{2.14}$$

This equality is known to be the compositeness condition in quantum field theory/11/.

As a result, the representation of the generating functional described mesons as the bound quark states is obtained in the form (2.13) under the condition (2.14). This representation is completely equivalent to (2.12) and underlies our model.

The baryons and other multiquark states can be constructed in the same manner. These states should arise from the terms L_n for $n \geq 3$ in (2.3). As in the case L_2, the gluon Green function $G^{a_1 \ldots a_n}_{\mu_1 \ldots \mu_n}(p_1, \ldots, p_n)$ is supposed to be analytical at points $p_j = 0$. The product of colourless n-quark states, which can be identified with the corresponding hadrons, is given off the L_n by using Firz transformation. The coefficients of expansions of n-point gluon Green functions at $p_j = 0$ define, first, the quark composition and quantum numbers of n-quark hadron states, and second, the hadron mass spectrum. However, this problem is unsolved at present because the n-point confined gluon Green functions are unknown yet.

We will use equivalence of the representations (2.12) and (2.13) with the auxiliary condition (2.14). We start from the mass spectrum and the quantum numbers of the observable hadrons and do some assumption about their quark composition. Then, the hadron-quark interaction Lagrangian can be written and the coupling constant can be defined from the compositeness condition.

This approach allows us to describe all the hadron interactions and to receive estimations of the expansion coefficients of the gluon Green functions from the mass equation (2.11).

3. THE QUARK CONFINEMENT MODEL

Now we will formulate the quark confinement model (QCM) which is expected to describe the low-energy physics. The model is based on the following assumptions.

1. The spectrum of hadrons with mass m_H and quantum numbers J^{PC} is taken from the experimental data. The quark structure of a hadron is defined by SU_3, C, P, T symmetry. The hadron-quark interaction Lagrangian is constructed in the simplest way without derivatives.

2. The hadron-hadron interactions are described by the functional (2.13). The quark-hadron coupling constants are defined by the compositeness condition (2.14).

3. The confinement ansatz. The quark confinement and ultraviolet convergence of the Feynman diagrams are provided by averaging over the gluon vacuum fields.

We start from the discussion of the hadronization. Let us consider mesons, baryons and multiquark states.

Mesons (two-quark states). The colourless two-quark currents with meson quantum numbers $Q = I^G J^{PC}$ can be chosen in the form

$$J_Q = \bar{q}(x) \Gamma_Q q(x), \tag{3.1}$$

where matrix Γ_Q provides necessary quantum numbers. In fact there are several such matrices, e.g. current $\bar{q}(x) \Gamma_Q (\overleftarrow{\partial} \overrightarrow{\partial}) q(x)$ has the same quantum numbers as (3.1).

The connection of a meson field M_Q with the quark currents can be written in the form

$$\mathcal{L}_Q = g_Q M_Q(x) \sum_i c_i J_Q^i(x). \tag{3.2}$$

The sum is over all two-quark currents with the same quantum numbers. The coefficients c_i satisfy the condition $\sum_i c_i^2 = 1$. We will use the quark currents with the lowest derivatives, which corresponds to the choice of the two-quark states with the lowest orbital momentum. It is the only state for the most cases.

Baryons (three-quark states). The colourless three-quark currents with baryon quantum numbers $B = J^P$ can be written in the form

$$J_B(x) = R_B\, q^a(x)\, q^b(x)\, q^c(x)\, \epsilon^{abc}, \tag{3.3}$$

where R_B is the straight product of the spinor and flavour matrices. The connection of the baryon field B_B with the quark currents can be written as

$$\mathcal{L}_B = g_B\, \bar{B}_B \sum_j c_j\, R_B^j\, q^a q^b q^c\, \epsilon^{abc} + h.c., \tag{3.4}$$

where g_B is the coupling constant. The sum is over all three-quark currents with the same quantum numbers. By analogy with mesons, the baryon interaction Lagrangian must be obtained by means of the expansion of the three-gluon Green function at small momenta. Since, it is unknown yet, the Lagrangian is constructed with the lowest degree of the derivatives.

The multiquark states are constructed by analogy with mesons and baryons. The main difficulty is that there are many ways to construct the multiquark currents with the given quantum numbers. Therefore, we sould not give any formulae here.

The hadron-hadron interactions are described by the generating functional by analogy with (2.13):

$$\begin{aligned} Z[J_M, J_B] = &\int \prod_Q \delta M_Q \int \prod_B \delta B_B \delta \bar{B}_B \int \prod_f \delta q \delta \bar{q} \int d\sigma_{vac} \times \\ &\times \exp\Big\{ i \int dx \Big[\tfrac{1}{2} M_Q (\Box - m_Q^2) M_Q + \bar{B}_B (i\hat{\partial} - m_B) B_B + \\ &+ \bar{q}_f (i\hat{\partial} - m_f - \hat{B}_{vac}) q_f + \mathcal{L}_Q + \mathcal{L}_B + M_Q J_Q + \bar{B}_B J_B + \bar{J}_B B_B \Big] - \\ &- i \iint dx_1 dx_2 \Big[\tfrac{1}{2} M_Q\, g_Q^2\, \Pi_Q^{reg} M_Q + \bar{B}_B\, g_B^2\, \Sigma_B^{reg} B_B \Big] \Big\}. \end{aligned} \tag{3.5}$$

The Lagrangians $\mathcal{L}_Q$ and $\mathcal{L}_B$ are defined by (3.2) and (3.4). The coupling constants are defined from the compositeness condition (2.14).

Using (3.5) one can get the following representation for the S -matrix:

$$\begin{aligned} S &= \int d\sigma_{vac}\, T \exp\Big\{ i \int dx \Big(\sum_Q \mathcal{L}_Q(x) + \sum_B \mathcal{L}_B(x) \Big) + \delta L \Big\}, \\ \delta L &= -i \iint dx_1 dx_2 \Big[\frac{1}{2} \sum_Q M_Q\, g_Q^2\, \Pi_Q^{reg} M_Q + \sum_B \bar{B}_B\, g_B^2\, \Sigma_B^{reg} B_B \Big]. \end{aligned} \tag{3.6}$$

Here, the time-ordered product is supposed to be the ordinary Vick's T-product of the hadron and quark fields with the quark propagator

$$\overbrace{q_f(x) \bar{q}_{f'}(x')} = \delta_{ff'} \frac{1}{i} \left(m_f + \hat{B}_{vac}(x) - i\hat{\partial} \right)^{-1} \delta(x - x'). \tag{3.7}$$

The quark fields must be equal to zero in (3.6) after the normal ordering, The measure $d\sigma_{vac}$ is the same as in (2.10).

Further, we will use the representation (3.6) of the scattering matrix.

4. THE QUARK CONFINEMENT HYPOTHESIS

The confinement hypothesis concerns the definition of an action of the vacuum gluon field $\hat{B}_{vac}(x)$ on the quark fields. This action is expected to provide quark confinement and to determine the dynamics of hadron interactions in the low-energy region. At present, these questions are the least known from the view point of rigorous results of the theory. Our confinement hypothesis consists in that in the low-energy region quarks turn into some quasiparticles which do not exist as usual particles.

The ansatz is to change the integration over $d\sigma_{vac}$ of the quark loops in (2.10) by the one-multiple integral

$$\int d\sigma_{vac}\, \mathrm{tr} \left[M(x_1) S(x_1, x_2 | B_{vac}) \cdots M(x_n) S(x_n, x_1 | B_{vac}) \right] \to$$

$$\to \int d\sigma_\lambda \, \mathrm{tr}\left[M(x_1)S_\lambda(x_1-x_2)\cdot\ldots\cdot M(x_n)S_\lambda(x_n-x_1)\right], \qquad (4.1)$$

where

$$S_\lambda(x_1-x_2)=\int\frac{d^4p}{(2\pi)^4 i}\cdot\frac{e^{-ip(x_1-x_2)}}{\lambda\Lambda_f+m_f-\hat{p}-i0}. \qquad (4.2)$$

The mass dimensional parameter Λ_f depends on flavour and defines the confinement region of quarks with flavour f .

Thus, the confinement ansatz (4.1) implies that we use the one-multiple integral which effectively takes into account the action of the vacuum gluon fields on the quarks instead of the functionals (2.10) and (3.5).

Further, we would like to solve two problems. First, to provide the quark confinement, i.e. to guarantee the absence of singularities corresponding to the quark production in the S -matrix elements. Second, to achieve the ultraviolet convergence of all integrals defining the S-matrix elements. This allows us to restrict the number of open model parameters.

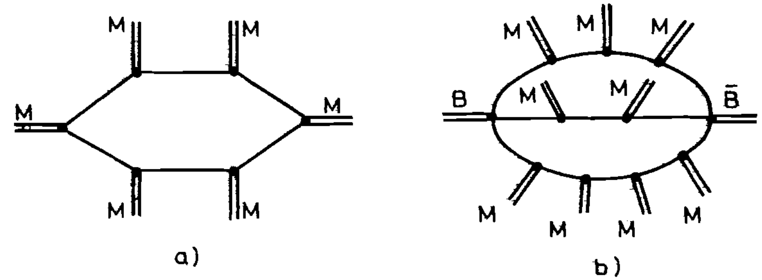
The measure $d\sigma_\lambda$ is defined along two lines: the algebraic structure of measure, which is connected with an action in the colour space, and the analytical one.

Algebraic Structure of Measure

After integrating over quark field in (3.5) and (3.6) the generating functional or S-matrix can be represented as a set of close colourless quark loops for the mesons and more complex structures for the baryons and multiquark states. First, the colourless quark loops are assumed to be averaged over the gluon measure independently

$$d\sigma_\lambda=\prod_{a=1}^{3}d\sigma_\lambda^a\,, \qquad \int d\sigma_\lambda^a=1\,.$$

For example, the meson-meson interactions (see, Fig. 1a) are described by the following integral:



F i g. 1. The quark diagrams describing the meson-meson (a) and meson-baryon interaction (b).

$$\int d\sigma_{vac}\,\mathrm{tr}\left[M(x_1)S(x_1,x_2|B_{vac})\cdot\ldots\cdot M(x_n)S(x_n,x_1|B_{vac})\right]\to$$

$$\to 3\int d\sigma_\lambda\,\mathrm{tr}\left[M(x_1)S_\lambda(x_1-x_2)\cdot\ldots\cdot M(x_n)S_\lambda(x_n-x_1)\right]. \qquad (4.3)$$

The meson-baryon amplitudes (see, Fig. 1b) are defined as

$$\int d\sigma_{vac}\,\epsilon^{abc}\,I^{aa'}S^{(k_1)}(x_1,x_2|B_{vac})I^{bb'}S^{(k_2)}(x_1,x_2|B_{vac})\times$$
$$\times I^{cc'}S^{(k_3)}(x_1,x_2|B_{vac})\,\epsilon^{a'b'c'}\to 6\prod_{j=1}^{3}S^{(k_j)}(x_1,x_2), \qquad (4.4)$$

where

$$S^{(k)}(x_1,x_2)=\int dy_1\cdot\ldots\cdot\int dy_k\int d\sigma_\lambda\,S_\lambda(x_1-y_1)M(y_1)\cdot\ldots\cdot M(y_k)S_\lambda(y_k-x_2). \qquad (4.5)$$

We will use the function (4.5) for any quark line with the given colour in more complex diagrams.

Analytical Structure of Measure

The confinement ansatz consists in the definition of the analytical structure of measure $d\sigma_\lambda$, i.e. in the definition of integral

$$\int \frac{d\sigma_\lambda}{\lambda \Lambda_f + m_f - \hat{p}} = \frac{1}{\Lambda_f} G\left(\frac{m_f - \hat{p}}{\Lambda_f}\right) = \tag{4.6}$$

$$= \frac{1}{\Lambda_f}\left\{ a\left(-\frac{p^2}{\Lambda_f^2}, \frac{m_f}{\Lambda_f}\right) + \frac{\hat{p}}{\Lambda_f} b\left(-\frac{p^2}{\Lambda_f}, \frac{m_f}{\Lambda_f}\right)\right\}$$

or

$$\int \frac{d\sigma_\lambda}{\lambda + z} = G(z). \tag{4.7}$$

The function $G(z)$ is called the confinement function; $G(z)$ is supposed to satisfy the following conditions:

1. $G(z)$ is a universal function, i.e. is independent of colour and flavour. In other words, the function $G(z)$ is unique for all quark diagrams defining the hadron-hadron interactions.

2. $G(z)$ is an entire analytical function on z-plane. This condition provides the quark confinement and unitarity of S-matrix on the hadron states.

3. $G(z)$ decreases faster than any degree of z in Euclidean region

$$\lim_{z^2 \to -\infty} (-z^2)^N |G(z)| = 0 \quad , \forall N > 0 .$$

All the quark diagrams converge due to this condition.

All the calculations are carried out in the Euclidean metric and then the analytical continuation to the physical momentum in the S-matrix elements is performed.

The S-matrix is constructed to be finite, unitary and macrocausal on the space of physical hadron states in each perturbative order/12/.

The shape of $G(z)$ in (4.7) is not yet obtained from any general QCD researches. Therefore, the choice of this shape is one of the model assumptions.

The function

$$G(z) = e^{(1-z)^2} \tag{4.8}$$

was used in refs./13/. In this case

$$a\left(-\frac{p^2}{\Lambda^2}, \frac{m}{\Lambda}\right) = \cos\left(\frac{2m}{\Lambda}\sqrt{-\frac{p^2}{\Lambda^2}}\right) e^{\frac{p^2}{\Lambda^2}},$$

$$b\left(-\frac{p^2}{\Lambda^2}, \frac{m}{\Lambda}\right) = \frac{\sin\left(\frac{2m}{\Lambda}\sqrt{-\frac{p^2}{\Lambda^2}}\right)}{\sqrt{-\frac{p^2}{\Lambda^2}}} e^{\frac{p^2}{\Lambda^2}},$$

where m and Λ are the open parameters defined by fitting of experimental data. The choice of the confinement function in the form (4.8) permited to describe the numerous phenomena of the meson and baryon physics at low energy from the unique point of view/4,13/.

However, matrix elements of physical processes have an exponential growth by energy due to (4.8). The auxiliary assumptions are demanded/13/ to extend the region of model application up to energy $\sim$ 1.5 - 2 GeV.

In the QCM we will proceed from the next assumptions under the choice of confinement function.

First, the functions $a(z^2, \xi)$ and $b(z^2, \xi)$ in (4.6) are considered to be independent since the $G(z)$ is an entire one and the form (4.6) is the most general representation of entire function.

Second, let us take the assumption concerning a way of breaking of the flavour group. At present, we are still far from understanding such questions as the SU_3 -breaking and quark mass splitting. Therefore, the quark masses are supposed to be different in the QCD Lagrangian (2.1). Then, the confinement is known to be the ansatzs (4.1) and (4.2) where defines the confinement region of a quark f . Generally speaking, the parameters m_f and Λ_f are independent at this stage. However, one can assume that Λ_f is proportional to m_f i.e. the relation

$$\xi = \frac{2m_f}{\Lambda_f} \tag{4.9}$$

is the same for all flavours. Then, the definition of the confinement function (4.6) is written as

$$\int \frac{d\sigma_\lambda}{\lambda\Lambda_f + m_f - \hat{p}} = \frac{1}{\Lambda_f}\left[a\left(-\frac{p^2}{\Lambda_f^2}\right) + \frac{\hat{p}}{\Lambda_f}\, b\left(-\frac{p^2}{\Lambda_f^2}\right)\right] \tag{4.10}$$

where

$$a(z^2) = \int \frac{d\sigma_\lambda\, \lambda}{\lambda^2 + z^2}, \quad b(z^2) = \int \frac{d\sigma_\lambda}{\lambda^2 + z^2}. \tag{4.11}$$

It is natural that the function $a + zb$ satisfies all the above-mentioned conditions. The suggestion (4.9) diminishes the number of the model dimensional parameters and the representations (4.10) and (4.11) are more convenient for the calculations.

Third, the functions $a(-k^2)$ and $b(-k^2)$ are assumed to be fast decreased both in the Euclidean direction and the Minkowski one

$$\lim_{k^2\to\pm\infty} |k^2|^N |a(-k^2)| = \lim_{k^2\to\pm\infty} |k^2|^N |b(-k^2)| = 0$$

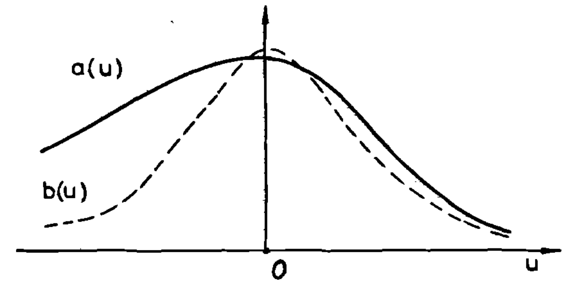
for any $N > 0$. It provides the absence of the exponential growth of matrix elements with energy.

In this paper we will use the simplest shape of $a(u)$ and $b(u)$

$$\begin{aligned} a(u) &= a_0 \exp\{-u^2 - 2a_1 u\}, \\ b(u) &= b_0 \exp\{-u^2 + 2b_1 u\}, \end{aligned} \tag{4.12}$$

where the parameters a_0, a_1, b_0, b_1 are defined by fitting over experimental data.

One has to remark that different shapes of $a(u)$ and $b(u)$ satisfying the general requirements were used for a description of the main meson decays. It is found that a behaviour of these functions must be similar as is shown in Fig.2 and results are stable as regards to change of the shapes close to ones in Fig. 2.



F i g. 2. The possible shapes of $a(u)$ and $b(u)$.

In the QCM there is a single dimensional parameter Λ_f defined the confinement region of the quark f. This parameter can be connected with the quark condensate:

$$\begin{aligned} \langle \bar{q}_f q_f \rangle &= -\,\mathrm{tr} \int \frac{d^4k}{(2\pi)^4 i} \int d\sigma_\lambda \frac{1}{\lambda\Lambda_f - \hat{k} - i0} = \\ &= -\Lambda_f^3 \frac{3}{(2\pi)^2} \int_0^\infty du\, u\, a(u). \end{aligned} \tag{4.13}$$

Finally, one can remark that the so-called virton-quark model[4] can be obtained on the basis of the above considerations if the measure $d\sigma_\lambda$ is defined as

$$\int d\sigma_{vac} \prod_{i,j\in\mathcal{D}} S(x_i, x_j | B_{vac}) \to \prod_{i,j\in\mathcal{D}} G(x_i - x_j),$$

$$G(x_i - x_j) = \int d\sigma_\lambda\, S_\lambda(x_i - x_j)$$

for any quark diagram $\mathcal{D}$ and the Fourier-transform of the quark propagator $\tilde{G}(\hat{p})$ is chosen as (4.8).

5. SCATTERING MATRIX IN QCM

The scattering matrix (3.6) describes all possible processes of the hadron interactions. There are hadron and quark field in the representation (3.6) but quarks in the form of the free asymptotic fields are absent due to the confinement ansatz.

The hadron-quark coupling constants are calculated from the compositeness condition (2.14) that is the strong coupling condition. Thus, the QCM describes strong interactions. Therefore, the perturbative theory over the coupling constant is not applicable. What is the way of the calculations? First of all, it is to be noted that the chain approximation is assumed to be valid for the compositeness condition [11]. Therefore, our calculations will also be based on the approximations which are connected with assuming up of the same class of diagrams for the hadron Green functions.

These approximations will be done in the framework of the $1/N_c$--expansion. Let us consider the meson-meson interactions. There are only quark loops in the representation of an S-matrix in the form of the Feynman diagrams. As follows from the compositeness condition (2.14) the effective strong coupling constant is

$$h_Q = 4N_c\left(\frac{g_Q^2}{4\pi}\right)^2, \qquad (5.1)$$

where $N_c = 3$ is the number of colours and 4 is the number of quark spinor indices. Therefore, the series of the perturbative theory can be represented over two parameters: h_Q and

$$\lambda = \frac{1}{4N_c} = \frac{1}{12}. \qquad (5.2)$$

Our approximations will be connected with the expansion over the parameter λ. Let us consider the first two approximations in detail.

First Approximation

It is called the one-loop approximation. Further, we will consider the meson Green function (pseudoscalar and scalar mesons) only. The meson Green function is defined by the one-loop quark diagrams in the chain approximation (see, Fig. 3a) corresponding to the zero degree of the parameter λ and is written in the form allowing for the mass renormalization

$$G_1(p^2) = \frac{1}{m^2-p^2+h[\Pi_1(p^2)-\Pi_1(m^2)]} \xrightarrow[p^2\to m^2]{} \frac{Z_1}{m^2-p^2}$$
$$Z_1 = \frac{1}{1-h\,\Pi_1'(m^2)}. \qquad (5.3)$$

Here the notation $h\,\Pi_1(p^2) = g^2\widetilde{\Pi}(p^2)$ is accepted with correspondence of the definitions (2.11) and (5.1). The meson mass operator $\Pi_1(p^2)$ in (5.3) corresponds to the one-loop quark diagram.

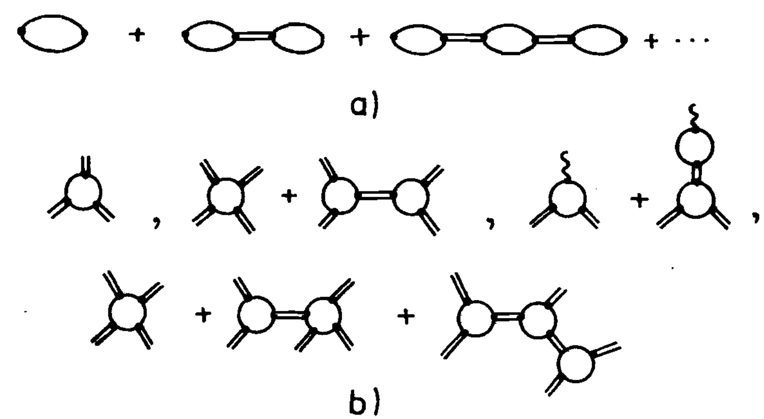


Fig. 3. The diagrams of the first (one-loop) approximation.

The renormalized coupling constant is defined by the relation

$$h_r = Z_1 h$$

whence

$$Z_1 = 1 + h_r\,\Pi_1'(m^2).$$

From the compositeness condition (2.14) we have

$$h_r = \frac{1}{(-\Pi_1'(m^2))}. \qquad (5.4)$$

One can obtain the following expression for the Green function:

$$h G_1(p^2) = \frac{Z_1 h}{Z_1(m^2-p^2) - Z_1 h[\Pi_1(m^2) - \Pi_1(p^2)]} = h_2 G_{12}(p^2), \tag{5.5}$$

$$G_{12}(p^2) = \frac{\Pi_1'(m^2)}{\Pi_1(m^2) - \Pi_1(p^2)} .$$

Let us discuss the analytical properties of this function. The mass operator $\Pi_1(p^2)$ is an entire analytical function on the p^2--plane. The Green function $G_{12}(p^2)$ in the one-loop approximation has a single pole on the real axis at $p^2 = m^2$ which defines meson mass. There are poles on the complex p^2-plane which correspond to the zeros of the function $\Pi_1(m^2) - \Pi_1(p^2)$. These poles are unphysical and connected with the λ -expansion used. It is essential that the function $G_{12}(p^2)$ tends to a constant as $p^2 \to -\infty$.

Thus, the one-loop approximation describes the behaviour of the hadron Green functions only near the mass shell, i.e. in the infrared limit.

Then note that the one-loop approximation in the representation (3.5) corresponds to a combination of the free meson Lagrangian with the one-loop counterterm δL

$$G_{12}^{-1}(p^2) = m^2 - p^2 + g^2 \tilde{\Pi}_{reg}(p^2) =$$

$$= m^2 - p^2 - \frac{1}{\tilde{\Pi}'(m^2)} [\tilde{\Pi}(p^2) - \tilde{\Pi}(m^2) - \tilde{\Pi}'(m^2)(p^2 - m^2)] =$$

$$= \frac{\tilde{\Pi}(m^2) - \tilde{\Pi}(p^2)}{\tilde{\Pi}'(m^2)} = \frac{\Pi_1(m^2) - \Pi_1(p^2)}{\Pi_1'(m^2)} .$$

The hadron amplitudes in the one-loop approximation are described by the one-loop quark diagrams which are the tree diagrams over the hadron lines. One can see the examples of such diagrams in Fig.3b. The hadron propagator in the one-loop approximation (5.5) corresponds to every internal hadron line.

The quark loops are shown in Fig. 3b to be the analytical functions over the external hadron momenta. Hence, there are no any threshold singularities. On the one hand, this corresponds to the complete quark confinement, and on the other hand, it shows that the one-loop approximation is applicable to the low-energy only, when the influence of possible intermediate processes can be neglected.

Thus, the one-loop approximation completely corresponds to the tree diagrams of the chiral theory. However, the difference is that the hadron-hadron vertices are the structureless points in the chiral theory, which corresponds to the local hadron interactions, while in the QCM they are described by the quark loops defining the hadron structure (form factors, slope parameters, etc.).

For completeness, let us obtain the vector field propagator in the one-loop approximation. The free vector propagator is given by

$$G_{0\mu\nu}(p) = \frac{1}{m_0^2 - p^2} \left[-g_{\mu\nu} + \frac{p_\mu p_\nu}{m_0^2} \right],$$

where m_0 is the bare mass of vector particles. The mass operator of the vector fields defined by a quark loop has the following shape due to the gauge invariance:

$$h \Pi_{\mu\nu}(p) = \left(g_{\mu\nu} - \frac{p_\mu p_\nu}{p^2} \right) h \Pi_V(p^2). \tag{5.6}$$

The vector field propagator $G_{\mu\nu}$ in the one-loop approximation is defined by the diagrams in Fig. 1a, i.e. it is the solution of the equation

$$G = G_0 + G_0 h \Pi G$$

and is written as

$$G_{\mu\nu} = \frac{-g_{\mu\nu} + \frac{p^\mu p^\nu}{p^2} \cdot \frac{p^2 - h\Pi_V(p^2)}{m^2 - h\Pi_V(m^2)}}{m^2 - p^2 + h[\Pi_V(p^2) - \Pi_V(m^2)]},$$

where

$$m^2 = m_0^2 + \delta m^2, \quad \delta m^2 = h \Pi_V(m^2).$$

The renormalization constant of the vector meson wave function is equal to

$$Z = \frac{1}{1 - h\,\Pi_V'(m^2)} = 1 + h_2\,\Pi_V'(m^2),$$

where the renormalized coupling constant is defined in a standard manner $h_2 = hZ$. If one assumes that $Z = 0$, then

$$h_2 = \frac{1}{(-\Pi_V'(m^2))} . \qquad (5.7)$$

Passing to the renormalized values in the relation

$$h\,G_{\mu\nu} = h_2\,G^{(2)}_{\mu\nu}$$

one can finally get

$$G^{(2)}_{\mu\nu} = \frac{\Pi_V'(m^2)}{\Pi_V(m^2) - \Pi_V(p^2)} \left[-g_{\mu\nu} + \frac{p_\mu p_\nu}{p^2} \cdot \frac{\Pi_V(p^2)}{\Pi_V(m^2)} \right] . \qquad (5.8)$$

Let us calculate the axial field propagator in the one-loop approximation. In this case the free axial propagator coincides with $G_{0\mu\nu}$. The axial mass operator is written as

$$h\,\Pi^A_{\mu\nu}(p) = h\left[g_{\mu\nu}\,\Pi_{tr}(p^2) - \frac{p_\mu p_\nu}{p^2}\,\Pi_l(p^2) \right].$$

Further calculations are carried out by analogy with the vector case. Finally, we obtain the following expression for the renormalized propagator

$$G^{(2)}_{\mu\nu}(p) = \frac{\Pi_{tr}'(m^2)}{\Pi_{tr}(m^2) - \Pi_{tr}(p^2)} \times \qquad (5.9)$$

$$\times \left[-g_{\mu\nu} + \frac{p_\mu p_\nu}{p^2} \cdot \frac{\Pi_l(p^2)}{\Pi_{tr}(m^2) - \Pi_{tr}(p^2) + \Pi_l(p^2)} \right]$$

where the coupling constant is equal to

$$h_2 = \frac{1}{(-\Pi_{tr}'(m^2))} .$$

Second Approximation

The linear in λ corrections to the hadron mass operator are defined by the diagrams shown in Fig. 4a,b. where the propagator in the one-loop approximation (5.5) puts in conformity to the internal meson lines

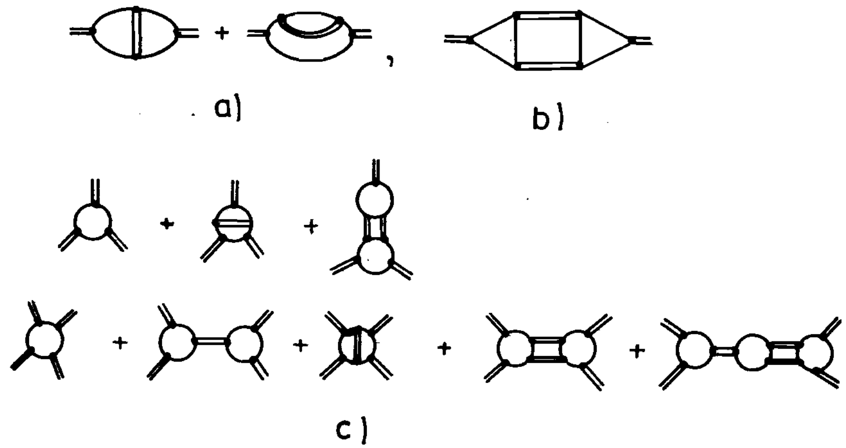


Fig. 4. The diagrams of the second approximation.

$$h_2\,\Pi_2(p^2) = h_2 \int d^4K\,\Gamma_4(p,K)\,\frac{1}{\Pi_1(K^2) - \Pi_1(m^2)} +$$

$$+ h_2 \iint d^4K_1\,d^4K_2\,\Gamma_3^2(K_1,K_2,p) \prod_{j=1}^{2} \frac{\delta^{(4)}(p-K_1-K_2)}{\Pi_1(K_j^2) - \Pi_1(m^2)} . \qquad (5.10)$$

Here the first term corresponds to the diagram of Fig. 4a and the second one to Fig. 4b. The functions $\Gamma_4(p,K)$ and $\Gamma_3(K_1,K_2,p)$ defining the quark loops decrease as $O(1/K^2)$ in the Euclidean metric as $K^2 \to \infty$ so that the hadron propagators in the one-loop approximation (5.5) tend to a constant. Therefore, the integrals in (5.10) diverge quadratically.

We will perform the renormalization of the hadron Green function so that the hadron mass and coupling constant are defined by the equations in the one-loop approximation. One can get

$$G_2^{(2)}(p^2)=\frac{\Pi_1'(m^2)}{\Pi_1(m^2)-\Pi_1(p^2)-\frac{\lambda}{\Pi_1'(m^2)}\Pi_2^{reg}(p^2)}, \tag{5.11}$$

$$\Pi_2^{reg}(p^2)=\Pi_2(p^2)-\Pi_2(m^2)-\Pi_2'(m^2)(p^2-m^2),$$

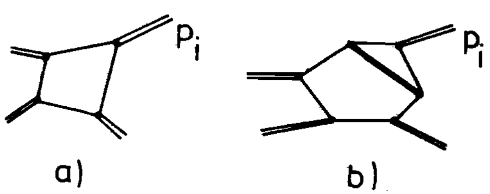
$$G_2^{(2)}(p^2)\to\frac{1}{m^2-p^2}\qquad(p^2\to m^2), \tag{5.12}$$

$$G_2^{(2)}(p^2)\to\frac{const}{(-p^2)\ln(-p^2)}\qquad(p^2\to-\infty). \tag{5.13}$$

The Green function of the second approximation (5.8) has a pole on the real axis at $p^2=m^2$ and the two-particle threshold singularities corresponding to the two-meson intermediate states. Moreover, there are unphysical singularities connected with λ -expansion.

In this approximation, the amplitudes of the hadron processes are described by the diagrams (for example, see Fig. 4b) where the internal meson lines are used to be the propagators (5.11). These amplitudes have the two-particle threshold singularities over the external hadron momenta. This approximation is applicable for higher energies than the first one.

The diagrams of the second approximation (see Fig. 5b) diverge logarithmically. The divergence is removed by the renormalization of the coupling constant in the diagram, Fig. 5a. Thus, we have



F i g. 5. The renormalization of the coupling constant in the second approximation.

$$T_2(\ldots,p_i^2,\ldots)\to T_2^{reg}(\ldots,p_i^2,\ldots)=T_2(\ldots,p_i^2,\ldots)-T_2(\ldots,m^2,\ldots) \tag{5.14}$$

The rules have been formulated to define completely the second approximation of λ -expansion in our model.

The subsequent approximations over the parameter λ can be constructed in an analogous way.

6. THE MAIN CHARACTERISTICS OF LIGHT MESONS AT LOW ENERGY

In the QCM hadron interactions are defined by the quark structure of hadrons and quark behaviour at large distances. It is essential that there is a possibility to calculate both the static values as lifetimes, magnetic moments, electromagnetic radii and the momentum dependences as slope parameters, form factors, etc. First of all, one has to choose the confinement functions a and b and to define the dimensional parameter Λ_f. We proceed with the determination of the model parameters by fitting the well established values being input parameters in the main phenomenological approaches. First, it is the constant of the pion weak decay f_π = 132 MeV which is the fundamental parameter in the chiral theory/14/. Second, it is the constant characterizing the transition $\rho\to\gamma$ which is the initial parameter in the vector dominance model (VDM)/15/. Third, these are constants defining the decays $\pi^0\to\gamma\gamma$ and $\omega\to\pi\gamma$. These decays are ordinary described by the Adler anomalies/16/. Finally, it is the constant of the strong decay $\rho\to\pi\pi$.

In the QCM these values are described by the corresponding quark diagrams. The examples of calculations are given in Appendix. All the calculations are carried out in the one-loop approximation.

The interaction Lagrangians of mesons M (P , S , V , A) with quarks are written as

$$\mathcal{L}_I=\sum_M\frac{g_M}{\sqrt{2}}\bar{q}_a\Gamma_M\lambda_M q_a M, \tag{6.1}$$

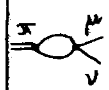
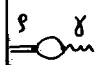
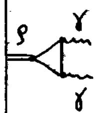
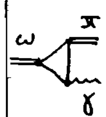
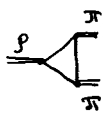
where Γ_M are the spin matrices ($i\gamma_5$, I , γ_μ , $\gamma_\mu\gamma_5$), and λ_M are the Gell-Mann isispin ones.

The coupling constants g_M are defined by the compositeness condition (2.14). The confinement function is of the shape (4.12). The best description of the experimental data is achieved for Λ = 480 MeV and

$$a_0 = 2.12 \qquad b_0 = 2$$
$$a_1 = 0.6 \qquad b_1 = 0.2 \tag{6.2}$$

in (4.12). The quark condensate (4.13) is equal to

T a b l e 1. The main characteristics of the light mesons described by the QCM

Process	Observable value	QCM	Expt./17/
	$f_\pi = \frac{\Lambda}{\pi}\frac{\sqrt{3}\,F_P(M_\pi^2)}{\sqrt{2F_{PP}(M_\pi^2)}}$	134 MeV	132 MeV
	$g_{\rho\gamma} = \frac{F_V(M_\rho^2)}{\pi\sqrt{8F_{VV}(M_\rho^2)}}$	0.18	0.20
	$g_{\pi\gamma\gamma} = \frac{F_{PVV}(M_\pi^2)}{\Lambda\pi\sqrt{3F_{PP}(M_\pi^2)}}$	0.27 GeV^{-1}	0.276 GeV^{-1}
	$g_{\omega\pi\gamma} = \frac{\sqrt{6}\,F_{PVV}(M_\omega^2)}{\Lambda\sqrt{F_{PP}(M_\pi^2)F_{VV}(M_\omega^2)}}$	2.25 GeV^{-1}	2.54 GeV^{-1}
	$g_{\rho\pi\pi} = \frac{\pi\sqrt{8}\,F_{VPP}(M_\rho^2)}{F_{PP}(M_\pi^2)\sqrt{F_{VV}(M_\rho^2)}}$	5.6	6.1

T a b l e 2. List of structure integrals

$F_{PP}(x) = \int_0^\infty du\, b(u) + \frac{x}{4}\int_0^1 du\, b(-\frac{ux}{4})\frac{(1-\frac{u}{2})}{\sqrt{1-u}}$
$F_{VV}(x) = \int_0^\infty du\, b(u) + \frac{x}{4}\int_0^1 du\, b(-\frac{ux}{4})\frac{(1-\frac{u}{2}+\frac{u^2}{4})}{\sqrt{1-u}}$
$F_V(x) = \int_0^\infty du\, b(u) + \frac{x}{4}\int_0^1 du\, b(-\frac{ux}{4})(1+\frac{u}{2})\sqrt{1-u}$
$F_P(x) = \int_0^\infty du\, a(u) + \frac{x}{4}\int_0^1 du\, a(-\frac{ux}{4})\sqrt{1-u}$
$F_{PVV}(x) = \frac{1}{4}\int_0^1 du\, a(-\frac{ux}{4})\, \ln\left(\frac{1+\sqrt{1-u}}{1-\sqrt{1-u}}\right)$
$F_{VPP}(x) = \int_0^\infty du\, b(u) + \frac{x}{4}\int_0^1 du\, b(-\frac{ux}{4})\sqrt{1-u}$

$$\langle \bar{u}u \rangle = -\frac{3\Lambda^3}{(2\pi)^2}\int_0^\infty dv\, v\, a(v) = -(182\ Mev)^3 \tag{6.3}$$

for the model parameters obtained.

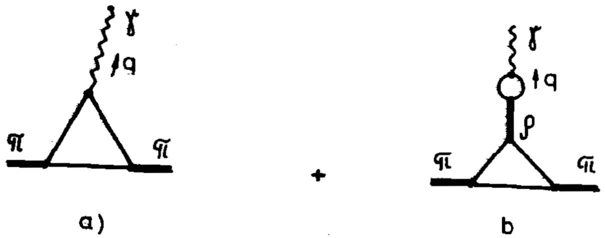
The fitted results are shown in Table 1. One can see that the fitting gives a good agreement of the theoretical values with the experimental data.

The next step in checking the main model assumption is the study of the momentum dependence of the physical matrix elements. The pion electromagnetic form factor is the best candidate for this.

Let us consider the behaviour of the pion elemctromagnetic form factor in the space-like region $q^2 = -Q^2 \le 0$, where q is the transfer momentum. The corresponding diagrams are shown in Figs.6a and b. The contributions of the triangle quark diagram in Fig. 6a and the resonance one in Fig. 6b are written as

$$F_\pi^{(a)}(q^2) = \frac{g_{\rho\pi\pi}(q^2)}{g_{\rho\pi\pi}(0)}\,, \tag{6.4}$$

$$F_\pi^{(b)}(q^2) = g_{\rho\pi\pi}(q^2)\,\mathcal{D}_\rho(q^2)\, q^2\, g_{\rho\gamma}(q^2)\,.$$



F i g. 6. The diagrams defining the pion electromagnetic form factor.

Here $\mathcal{D}_\rho(q^2)$ is the ρ -meson propagator (5.8). The expressions for $g_{\rho\gamma}(q^2)$ and $g_{\rho\pi\pi}(q^2)$ are represented in Tables 1 and 2. By using these formula we have

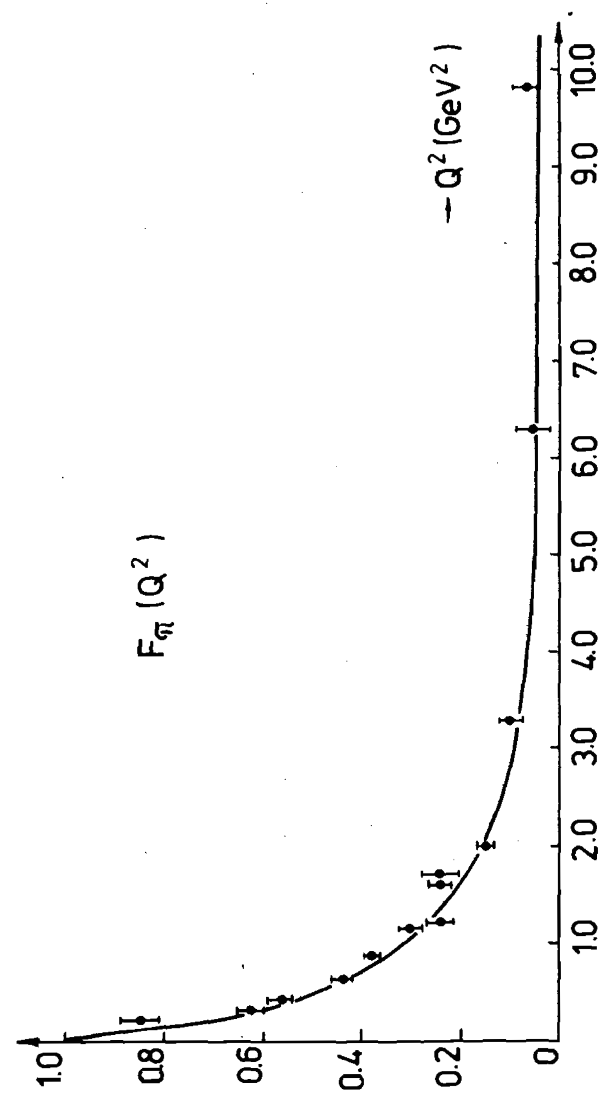
$$F_\pi(q^2) = F_\pi^{(a)}(q^2) + F_\pi^{(b)}(q^2) =$$

$$= \frac{F_{VPP}(q^2/\Lambda^2)}{F_{VPP}(0)} + \frac{h_\pi}{2}\cdot\frac{q^2\, F_{VPP}(q^2/\Lambda^2)\, F_V(q^2/\Lambda^2)}{m_\rho^2 F_V(m_\rho^2/\Lambda^2) - q^2 F_V(q^2/\Lambda^2)} = \tag{6.5}$$

$$= \frac{F_{VPP}(q^2/\Lambda^2)}{F_{VPP}(0)}\cdot\frac{m_\rho^2\, F_V(m_\rho^2/\Lambda^2)}{m_\rho^2 F_V(m_\rho^2/\Lambda^2) - q^2 F_V(q^2/\Lambda^2)}\,.$$

One can see, this expression coincides with corresponding one predicted by the VDM when the quark loops do not depend on momenta.

It is interesting to compare the contributions of the diagrams in Fig. 6a and Fig. 6b to the pion electromagnetic radius. We have



F i g. 7. The behaviour of the pion form factor in the space-like region.

$$\langle r_\pi^2\rangle^{(a)} = 6\,\frac{g'_{\rho\pi\pi}(0)}{g_{\rho\pi\pi}(0)} = \frac{6}{m_\rho^2}\left[\frac{m_\rho^2}{\Lambda^2\, 3\int_0^\infty du\, b(u)}\right] = 0.15\ fm^2,$$

$$\langle r_\pi^2\rangle^{(b)} = \frac{6}{m_\rho^2}\cdot\frac{F_V(0)}{F_V(m_\rho^2/\Lambda^2)} = 0.28\ fm^2, \qquad (6.6)$$

$$\langle r_\pi^2\rangle = \langle r_\pi^2\rangle^{(a)} + \langle r_\pi^2\rangle^{(b)} = 0.43\ fm^2.$$

One can see, the resonance diagram gives a smaller value than the VDM prediction $\langle r_\pi^2\rangle_{VDM}$ = 0.39 fm^2 due to the inclusion of the pion quark structure but the contribution of the triangle diagram increases the total value.

The behaviour of the pion form factor in the space-like region $q^2:\ 0 \le Q^2 = -q^2 \le 10\ GeV^2$ is shown in Fig. 7. One can see, the experimental data are described by our model quite accurately. It is to be remarked that the diagram in Fig. 6b, dominates for the large momenta $Q^2 \gtrsim 2$ GeV^2.

APPENDIX A: The Calculation Technique

Let us consider the two- and three-point quark loops to demonstrate the calculation technique of the matrix elements in the QCM.

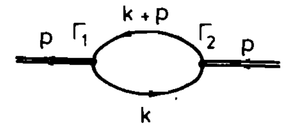
The two-point quark loop is shown in Fig. 8. The following integral:

$$\Pi_{12}(p) = 3g_1 g_2 \int\frac{d^4k}{(2\pi)^4 i}\int d\sigma_\lambda\, tr\left[\Gamma_1\frac{1}{\lambda\Lambda_1 - \hat{k} - \hat{p}}\Gamma_2\frac{1}{\lambda\Lambda_2 - \hat{k}}\right] \qquad (A.1)$$

corresponds to this diagram. Here, Γ_1 and Γ_2 are γ-matrices defining the interaction of a meson with quarks, Λ_1 and Λ_2 are the parameters describing the quark confinement region and g_1 and g_2 are the coupling constants.

By using the Feynman α-parametrization one can obtain

$$\Pi_{12}(p) = \frac{3g_1g_2}{(2\pi)^2}\int_0^1 d\alpha\int_0^\infty du\int d\sigma_\lambda\,\frac{R(u,\alpha,p)}{\left[\lambda^2 + u - \alpha(1-\alpha)\frac{p^2}{\Lambda^2(\alpha)}\right]}\,. \qquad (A.2)$$



F i g. 8. The two-point quark loop.

Here

$$R(u,\alpha,p) =$$
$$= \frac{1}{4}tr\left[\lambda^2\Lambda_1\Lambda_2\Gamma_1\Gamma_2 - \alpha(1-\alpha)\Gamma_1\hat{p}\Gamma_2\hat{p} - \frac{u}{4}\Lambda^2(\alpha)\Gamma_1\gamma^\alpha\Gamma_2\gamma^\alpha\right] +$$
$$+\lambda\frac{1}{4}tr\left[(1-\alpha)\Lambda_2\Gamma_1\hat{p}\Gamma_2 - \alpha\Gamma_1\Gamma_2\hat{p}\right];$$
$$\Lambda^2(\alpha) = \alpha\Lambda_1^2 + (1-\alpha)\Lambda_2^2 = \Lambda^2\left[1+\Delta(1-2\alpha)\right],$$
$$\Lambda^2 = \frac{\Lambda_1^2+\Lambda_2^2}{2}\ ,\quad \Delta = \frac{\Lambda_2^2-\Lambda_1^2}{\Lambda_2^2+\Lambda_1^2}.$$

Recalling the definitions (4.11), we have

$$\int_0^\infty du\int d\sigma_\lambda\frac{\lambda^2}{[\lambda^2+Q]^2} = -\int_0^\infty du\, Q\, b(Q),$$
$$\int_0^\infty du\int d\sigma_\lambda\frac{1}{[\lambda^2+Q]^2} = \int_0^\infty du\, b(Q), \qquad (A.3)$$
$$\int_0^\infty du\, u^2\int d\sigma_\lambda\frac{1}{[\lambda^2+Q]^2} = 2\int_0^\infty du\, u\, b(Q)$$

where $$Q = u - \alpha(1-\alpha)\frac{p^2}{\Lambda^2(\alpha)}\,.$$

The following equality is used

$$\int_0^1 d\alpha\, R(\alpha) \int_0^\infty du\, F\Big(u-\alpha(1-\alpha)\frac{p^2}{\Lambda^2(\alpha)}\Big) =$$

$$= \int_0^1 d\alpha R(\alpha) \int_0^\infty du\, F(u) + s\int_0^{u_\Delta} du\, F(-us) \int_{\alpha_-(u)}^{\alpha_+(u)} d\alpha\, R(\alpha), \qquad (A.4)$$

$$s = \frac{p^2}{4\Lambda^2}, \quad u_\Delta = \frac{2}{1+\sqrt{1-\Delta^2}}, \quad \alpha_\pm(u) = \frac{1}{2}\Big\{1+\frac{\Delta u}{2} \pm \sqrt{1-u+\Big(\frac{\Delta u}{2}\Big)^2}\Big\}.$$

Using formula (A.3) and (A.4) in the expression (A.2), one can obtain

$$\Pi_{12}(p) = -\Lambda^2 \frac{3g_1g_2}{(2\pi)^2}\cdot\frac{1}{4}\,\mathrm{tr}\, W_2,$$

$$W_2 = \Big[\Gamma_1\Gamma_2\sqrt{1-\Delta^2} + \frac{1}{2}\Gamma_1\gamma^6\Gamma_2\gamma^6\Big]\Big[b_1 - s^2\int_0^{u_\Delta} du\,u\, b(-us)\sqrt{1-u+\Big(\frac{\Delta u}{2}\Big)^2}\,\Big] +$$

$$+ \frac{\Delta^2}{4}\Gamma_1\gamma^6\Gamma_2\gamma^6\cdot s^2\int_0^{u_\Delta} du\,u^2\, b(-us)\sqrt{1-u+\Big(\frac{\Delta u}{2}\Big)^2} +$$

$$+\frac{1}{6}\Big[\Gamma_1\frac{\hat p}{\Lambda}\Gamma_2\frac{\hat p}{\Lambda} + 2s\Gamma_1\gamma^6\Gamma_2\gamma^6\Big]\Big[b_0 + s\int_0^{u_\Delta} du\, b(-us)\sqrt{1-u+\Big(\frac{\Delta u}{2}\Big)^2}\Big(1+\frac{u}{2}-\frac{\Delta^2 u^2}{2}\Big)\Big] +$$

$$+\frac{1}{2}\Big[\sqrt{1-\Delta}\,\Gamma_1\Gamma_2\frac{\hat p}{\Lambda} - \sqrt{1+\Delta}\,\Gamma_1\frac{\hat p}{\Lambda}\Gamma_2\Big]\Big[a_0 + s\int_0^{u_\Delta} du\, a(-us)\sqrt{1-u+\Big(\frac{\Delta u}{2}\Big)^2}\,\Big] +$$

$$+\frac{\Delta}{4}\Big[\sqrt{1-\Delta}\,\Gamma_1\Gamma_2\frac{\hat p}{\Lambda} + \sqrt{1+\Delta}\,\Gamma_1\frac{\hat p}{\Lambda}\Gamma_2\Big]\cdot s\int_0^{u_\Delta} du\,u\, a(-us)\sqrt{1-u+\Big(\frac{\Delta u}{2}\Big)^2}, \qquad (A.5)$$

$$a_0 = \int_0^\infty du\, a(u), \quad b_0 = \int_0^\infty du\, b(u), \quad b_1 = \int_0^\infty du\,u\, b(u).$$

When $\Lambda_1=\Lambda_2=\Lambda$ ($\Delta = 0$) the formula (A.5) **is simplified**

$$\Pi_{12}(p) = -\Lambda^2\frac{3g_1g_2}{(2\pi)^2}\,\frac{1}{4}\,\mathrm{tr}\Big\{\Big[\Gamma_1\Gamma_2 + \frac{1}{2}\Gamma_1\gamma^6\Gamma_2\gamma^6\Big]B_1(s) +$$

$$+\frac{1}{6}\Big[\Gamma_1\frac{\hat p}{\Lambda}\Gamma_2\frac{\hat p}{\Lambda} + 2s\Gamma_1\gamma^6\Gamma_2\gamma^6\Big]B_0(s) + \qquad (A.6)$$

$$+\frac{1}{2}(\Gamma_1\Gamma_2 - \Gamma_2\Gamma_1)\frac{\hat p}{\Lambda}A_0(s).$$

Here

$$B_1(s) = b_1 + s\int_0^1 du\,(-us)\,b(-us)\sqrt{1-u},$$

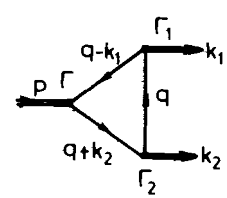
$$B_0(s) = b_0 + s\int_0^1 du\, b(-us)\sqrt{1-u}\,\Big(1+\frac{u}{2}\Big),$$

$$A_0(s) = a_0 + s\int_0^1 du\, a(-us)\sqrt{1-u}. \qquad (A.7)$$

Let us calculate the integral corresponding to the three-point quark loop (see, Fig. 9) when $K_1^2=K_2^2 = 0$ and $\Lambda_1=\Lambda_2=\Lambda_3=\Lambda$. The integral is written as

$$T(p,K_1,K_2) = 3g g_1 g_2 \int\frac{d^4q}{(2\pi)^4 i}\int d\sigma_\lambda \times$$

$$\times\,\mathrm{tr}\Big[\Gamma\frac{1}{\lambda\Lambda-\hat q+\hat K_1}\Gamma_1\frac{1}{\lambda\Lambda-\hat q}\Gamma_2\frac{1}{\lambda\Lambda-\hat q-\hat K_2}\Big]. \qquad (A.8)$$



F i g. 9. The three-point quark loop.

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Received by Publishing Department
on January 15, 1988.

Ефимов Г.В., Иванов М.А. E2-88-37
Конфайнмент и кварковая структура
легких адронов

Предлагается модель, называемая моделью конфайнмированных кварков /МКК/, предназначенная для описания низкоэнергетической физики легких адронов /мезонов и барионов/. Модель основана на двух гипотезах. Во-первых, исходя из идеи, что вакуум глюонных полей КХД обеспечивает конфайнмент кварков и любых цветных состояний, делается предположение о механизме усреднения промежуточных кварковых состояний по глюонному вакууму. Во-вторых, адроны рассматриваются как коллективные переменные, соответствующие бесцветным кварковым токам. С единой точки зрения дано описание сильных электромагнитных и слабых взаимодействий мезонов и барионов при низких энергиях.

Работа выполнена в Лаборатории теоретической физики ОИЯИ.

Препринт Объединенного института ядерных исследований. Дубна 1988

Efimov G.V., Ivanov M.A. E2-88-37
Confinement and Quark Structure
of Light Hadrons

We present a quark confinement model (QCM) for the description of the low-energy physics of light hadrons (mesons and baryons). The model is based on two hypotheses. First, the quark confinement is realized as averaging over vacuum gluon fields which are believed to provide the confinement of any colour objects. Second, hadrons are treated as collective colourless excitations of quark-gluon interactions. The description of strong, electromagnetic and weak interactions of mesons and baryons at the low energy is given from a unique point of view.

The investigation has been performed at the Laboratory of Theoretical Physics, JINR.

Preprint of the Joint Institute for Nuclear Research. Dubna 1988