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QUARK MODEL AND QCD

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ABSTRACT

The following topics are discussed:

1. QCD in the domain of soft processes
2. Phenomenological Lagrangian for soft processes and exotic mesons
3. Spectroscopy of low lying hadrons:
   (i) mesons
   (ii) baryons
   (iii) mesons with heavy quarks (c,b)
4. Confinement forces
5. Spectral integration over quark masses.

В.В. Анисович: Кварковая модель и КХД. KFKI-1989-32/A

АННОТАЦИЯ

Рассматриваются следующие темы:

1. КХД в области мягких процессов
2. Феноменологический лагранжиан для мягких процессов и экзотичных мезонов
3. Спектроскопия низколежащих адронов
   а) мезонов
   б) барионов
   в) мезонов с тяжелыми кварками (c, b)
4. Силы конфайнмента
5. Спектральные интегралы по массам кварков.

Anyiszovics V.V.: Kvark modell és QCD. KFKI 1989 32/A

KIVONAT

A következő témákat tárgyaljuk:

1. Kvantumszíndinamika lágy folyamatokban
2. Fenomenologikus Lagrangian lágy folyamatok és egzotikus mezonok esetére
3. Spektroszkópia:
   (i) mezonok
   (ii) barionok
   (iii) nehéz kvarkokat (c,b) tartalmazó mezonok
4. Confinement erők
5. Spektrálintegrálás kvark tömegekre

1. QCD IN THE DOMAIN OF SOFT PROCESSES

For soft physics it is possible to use two alternative languages: the language of hadrons and the language of quarks and gluons. Using the language of the quarks and gluons for description of the soft hadron physics it is necessary to take into account two characteristic phenomena which prevent one from usage of QCD Langragian in the straightforward way, ~~These~~ are

(i) chiral symmetry breaking, and

(ii) confinement of colour particles.

Chiral symmetry breaking reveals itself in hadron spectroscopy This symmetry breaking is realized in masses of the light quarks. In the region of soft physics the light quarks are much heavier than QCD quarks.

Table 1.

QCD quarks	Dressed quarks of the soft physics
$m_u \simeq$ 4 MeV	$m_u \simeq m_d \sim$ 300-400 MeV
$m_d \simeq$ 7 MeV	$m_s \sim$ 450-500 MeV
$m_s \simeq$ 150MeV	

Confinement reveals itself as the absence of free colour particles. Another consequence of the confinement is the massiveness of the effective gluons (e.g. gluons in the soft region). Let us look at an example: meson scattering amplitude with the t-channel gluon exchanges (fig. 1a).

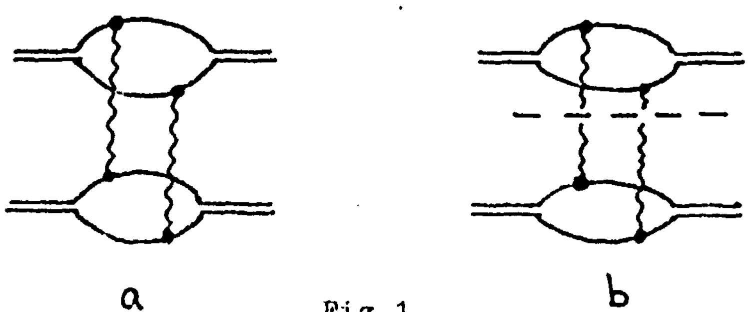

Fig.1.

Zero gluon mass leads to the t-channel singularity at t=0 and to $disc_t\ A \neq 0$ at $t > 0$ for the scattering amplitude A (the t-channel cutting of this diagram is shown in fig. 1b).

Due to confinement the amplitude A has t-channel singularities only at $t \geq 4m_\pi^2$. So the confinement requires that the amplitude discontinuity at $0 < t < 4m^2$ vanishes: $disc_t A=0$.

Most of the estimations of the effective gluon mass gives $500\ MeV \lesssim m_G \lesssim 1000\ MeV$. The massive gluon should be some complicated system of the QCD-gluons. The rules of the $1/N_c$--expansion (N_c is the number of quark colours) argue in favour of such a treatment of the effective gluon.

Let us consider the self-energy part of meson due to the interaction of the QCD-quarks and gluons. According to the $1/N_c$-rules the main contribution is given by the planar diagrams of the type fig. 2. Planar diagrams might be divided into parts with one $q\bar{q}$ pair in the intermediate state. This allows us to separate the blocks (fig. 3a) where q and $\bar{q}$ interact via gluon exchanges (denote this block as V).

One can assume that a sort of dual expansion is valid for gluon block V and it can be represented as a set of graphs with exchanges by some effective mesons (see fig. 3b):

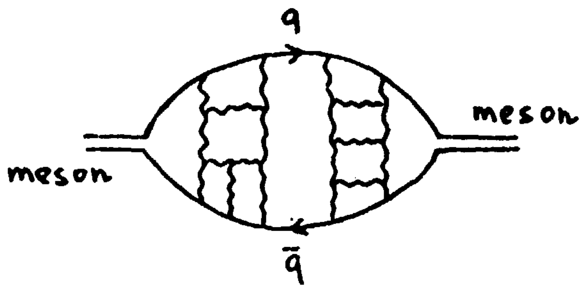

Fig. 2.

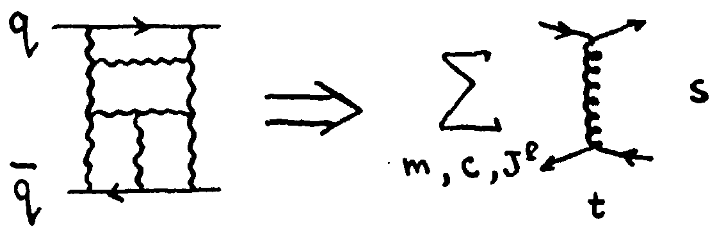

Fig. 3 (c-colour index)

$$V = \sum_{m,c,J^P} B_{m,c,J^P}(s) \frac{1}{m^2(c,J^P) - t} . \qquad (1)$$

Restricting our considerations by the region of low s (the region where the quark model is assumed to be valid) the factors B_{m,c,J^P} can be considered as independent of s. In this case it is enough to summarize over the lowest values J^P only.

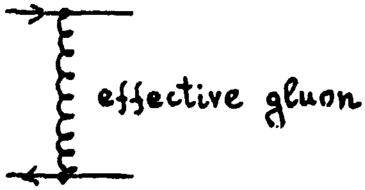

Fig. 4.

The calculation of hadron spectra in the framework of the quark model emphasizes the importance of the effective gluon exchange. It leads to the colour-magnetic forces which result in the mass splitting in the $SU(6)_{sp+fl}$ -multiplets.

The η-η' splitting and small value of pion mass indicate the existence of interactions in the white channels with $J^P=0^-$.

The role of other interactions is not clear so far.

2. PHENOMENOLOGICAL LAGRANGIAN AND EXOTIC MESONS

It seems rather perspective to construct a phenomenological Lagrangian for soft processes. Simple realization of such a renormalizable Lagrangian may be obtained when three types of fields are introduced: quark (q), gluon (A), and Higgs field (ϕ). Let us see a toy-example of such Lagrangian

$$\mathcal{L} = -\frac{1}{4} G_{\mu\nu}^2 + \bar{q}(i\hat{\mathcal{D}} - m)q - \frac{1}{2}\mathrm{Tr}\left((\mathcal{D}_\mu \phi)^+ (\mathcal{D}_\mu \phi)\right) + V(\phi, \phi^+) \qquad (2)$$

$$G_{\mu\nu}^a = \partial_\mu A_\nu^a - \partial_\nu A_\mu^a + g f^{abc} A_\mu^b A_\nu^c$$

$$\mathcal{D}_\mu = \partial_\mu - \frac{i}{2} g \lambda^a A_\mu^a$$

$\phi^{k\ell}$ — $SU(3)_g$-triplet (index k), $SU(3)_l$-triplet (index ℓ)

We will suppose that the global colour symmetry is not destroyed in the soft-physics region. A simple mechanism for obtaining the global colour symmetry and massive gluons is to introduce a scalar Higgs field which is a triplet of both the local colour group $SU(3)_l$ and some global $SU(3)_g$.

Potential V provides the spontaneous breaking of the local symmetry. Here it is enough to restrict oneself to the fourth degree of ϕ :

$$V = a\,\mathrm{Tr}(\phi^+\phi) + b\,\mathrm{Tr}(\phi^+\phi\phi^+\phi) + c(\mathrm{Tr}(\phi^+\phi))^2 + d(\det\phi^+ + \det\phi) \qquad (3)$$

If $b > 0$, $c > 0$, $a < 0$ and $d < 0$, the symmetry $SU(3)_l \times SU(3)_g$ is broken and the Lagrangian turns out to be invariant under global $SU(3)$-symmetry (global colour symmetry).

For Lagrangian with broken local symmetry we have following particles:

(i) massive gluons $G(J^P = 1^-)$ - octet of the global colour symmetry

(ii) octet of colour Higgs particles $H(0^+)$

(iii)two white Higgs particles $H_1(0^+)$, $H_1'(0^+)$.

Varying the parameters a, b, c, d, it is easy to obtain the masses of G, H, H_1 and H_1' in the range 1-2 GeV:

$$M_G, M_H, M_{H1}(0^+), M_{H'1}(0^+) \sim 1 - 2 \text{ GeV}. \quad (4)$$

It is reasonable to suppose that the masses of all composite particles with L = 0

$$\begin{array}{lll} (GG) & J^P = 0^+, 1^+, 2^+ & \\ (GH) & 1^- & (5) \\ (HH) & 0^+ & \end{array}$$

are of the same order of value 1-2 GeV. (The composite particles with colour constituents are listed here: the rules of $1/N_c$-expansion show that the colour constituents are likely to bind into composite systems).

We have here four white particles with $J^P = 0^+$. This corresponds to the situation observed experimentally: there is larger number of scalars among exotic mesons in the mass region 1-2 GeV.

So this toy-Lagrangian gives us not only massive gluons and global colour symmetry but scalar meson 'surplus' which is in a qualitative agreement with data (see results of the GAMS-experiment). However, in order to receive quantitative mass values we should make rather cumbersome calculation. It is seen from consideration of the spectroscopy of low-lying mesons which is discussed below.

3. SPECTROSCOPY OF LOW-LYING MESONS

Let us consider scattering amplitudes of constituent (dressed) quarks of three flavour (u,d,s): the poles of these amplitudes determine the masses of mesons. Constituent quark is colour triplet and quark amplitude obeys the global colour symmetry.

We assume the blocks with gluon exchanges are short-range ones (the masses of the effective mesons in dual expansion are large enough). For the low-energy quark processes we reduce these blocks to the point-like four-quark interaction (fig.5).

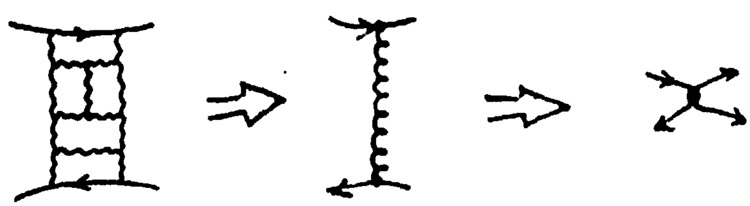

Fig. 5.

Four-quark interactions of two types are considered. Gluon-type exchange:

$$g_V\ (\bar{q}\ \vec{\lambda}\ \gamma_\mu\ q)^2. \tag{6}$$

Interaction due to the white exchanges:

$$g_I\left[(\bar{q}\gamma_5 q)^2 + (\bar{q}q)^2 - (\bar{q}\gamma_5\vec{\tau}q)^2 - (\bar{q}\vec{\tau}q)^2\right] + \\ + g_s\left[(\bar{s}\gamma_5 s)^2 + (\bar{s}s)^2\right] \tag{7}$$

Here q = u,d and $\vec{\tau}$ is isospin Pauli matrix. This interaction can be considered as induced by instantons.

We construct the low-energy amplitudes $qq \to qq$ and $q\bar{q} \to q\bar{q}$ with the help of some bootstrap iteration procedure considering four-fermion constants (g_I, g_s, g_v) as parameters.

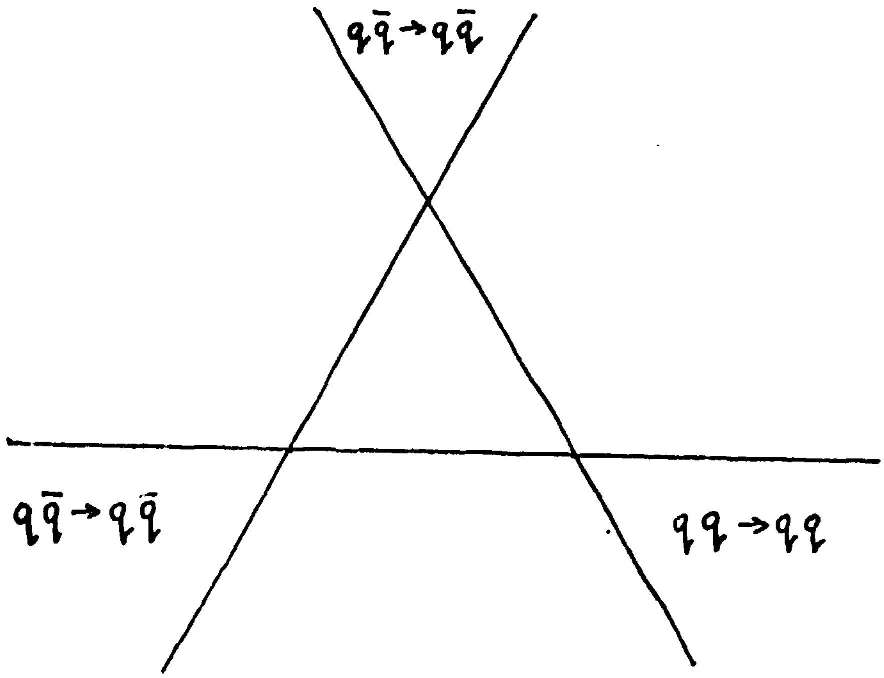

Fig. 6.

The scheme of the iteration bootstrap procedure is a follows. Partial amplitudes in all three channels (Mandelstam plane is shown in fig. 6) are calculated through the dispersion N/D-method, the above point-like interaction being the N-function (zero approximation) - see fig. 7. For regularization of the dispersion integrals we use a cut-off parameter Λ.

Fig. 7.

The amplitude of the first approximation is obtained when the zero approximation amplitude is taken as an N-function: the forces of the first approximation are determined by the exchange diagrams of the zero approximation (fig. 8) and so on. The iteration procedure converges very rapidly due to $1/N_c$-expansion (first approximation gives results which do not change practically in the following iterations).

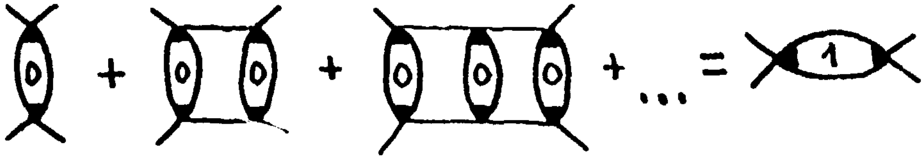

Fig. 8.

Strange quark violates the flavour $SU(3)_f$-symmetry. In order to avoid introducing additional parameters we use the dimentionless ones

$$\gamma_i = \frac{(m_a+m_b)^2}{4\pi^2} g_i \quad , \quad \lambda = \frac{4\Lambda}{(m_a+m_b)^2} \tag{8}$$

which are common for all processes. Here m_a, m_b are quark masses and $i = v, I, s$.

We obtain the following masses of mesons of the low-lying multiplets (0^-, 1^-, 0^+) /1/.

0^-	m (GeV)	1^-	m (GeV)	0^+	m (GeV)
π (0.14)	<u>0.14</u>	ρ (0.77)	<u>0.77</u>	δ (0.78)	0.78
η (0.55)	0.48	ω (0.78)	0.77	S*(0.98)	0.87
η'(0.96)	<u>0.96</u>	φ (1.02)	1.00	ε (1.30)	1.16
K(0.50)	<u>0.50</u>	K*(0.89)	0.89	$\varkappa$ (1.35)	0.88
$\vartheta(0^-)$ = -30°		$\vartheta(1^-)$ = 28°		$\vartheta(0^+)$ = -80°	

Here m_π, $m_{\eta'}$, m_K, m_ρ are used for determination of calculation parameters.

Mixing η-η' angle in the linear mass formula of the quar model is $\vartheta_\ell(0^-) = -23°$ (our calculation gives $\vartheta(0^-) \simeq \vartheta_\ell(0^-)$).

The mixing angle $\vartheta(1^-)$ corresponds to the fact φ-meson is mainly built of s-quarks:

$$\varphi \simeq s\bar{s} + (\text{small admixture of non-strange quarks}) \tag{9}$$

Our calculation indicates an important role of interaction of instanton-induced type. Such an interaction is necessary to obtain both the pion mass and η-η' mass splitting. However the relative contribution of the instanton-induced interaction (g_I) is less than that of gluon-type exchange (g_V). This ratio is of the order of

$$g_I/g_V \simeq 1/4 \qquad (10)$$

It indicates that the gluon-type exchanges play the main role in the formation of the pion.

For the 0^+ multiplet the discrepancy between calculated and observed masses is more than for 0^- and 1^-. It is possible that this is due to the admixture of the high-excited states ($qq\bar{q}\bar{q}$ or glueballs), which are not taken into account.

Calculation gives a bound state in the gluon channel:

$$M_G \simeq 0.67 \text{ GeV} \quad \text{- constituent gluon} \qquad (11)$$

The bound state of the constituent gluon arises because of the diagrams with quark loops (fig. 9).

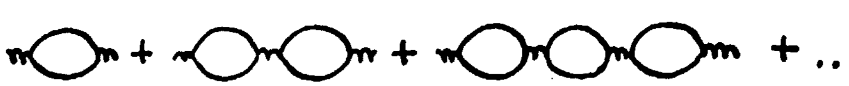

Fig. 9.

The calculated mass of the constituent gluon should be near the mass of the ρ-meson. It is a consequence of the rules of $1/N_c$-expansion (if $N_f = N_c$) and of the short-range type of the gluonic exchanges or the large value of the mass of the bare constituent gluon in above-mentioned phenomenological Lagrangian

$$M_G \text{ (bare)} > 1 \text{ GeV}. \qquad (12)$$

The value of the constituent gluon mass obtained here (~700 MeV) seems to be rather reasonable:

(i) this mass value is required by hard-process phenomenology

(ii) this mass being doubled is in the region of the masses of the most probable candidates for glueballs.

In our calculation there are diquark bound states with $J^P = 0^+$:

$$m_{ud}(0^+) \simeq 0.72 \text{ GeV} \quad \text{- isoscalar diquark}$$

$$m_{sd}(0^+) = m_{su}(0^+) \simeq 0.86 \text{ GeV - strange diquark} \qquad (13)$$

Bootstrap procedure gives us N-functions for qq- and $q\bar{q}$-interactions. These N-functions make it possible to calculate the form factors of the dressed quarks (fig. 10) and meson form factors (fig. 11).

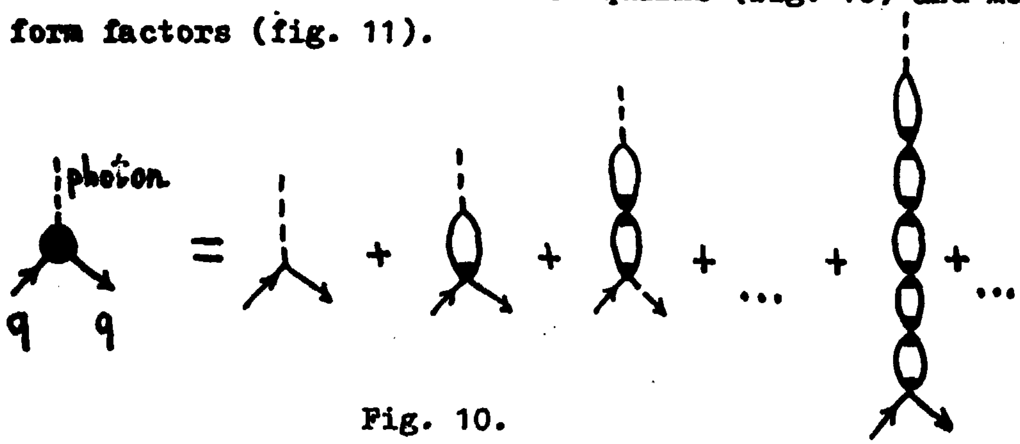

Fig. 10.

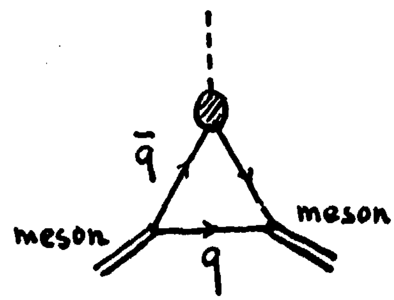

Fig. 11.

The results are as follows

$r_u \simeq 0.29$ fm, $r_d \simeq 0.26$ fm, $r_s \simeq 0.20$ fm, (14)

$R_\pi \simeq 0.78$ fm, $R_k \simeq 0.74$ fm (experimental value $R_\pi \simeq 0.62$ fm)

4. LOW-LYING BARYONS ($J^P = 1/2^+, 3/2^+$)

If we know qq-interaction, it is possible to consider baryon with $J^P = 1/2^+, 3/2^+$. This has been made using Faddeev equation (i.e. qq-interaction was considered in the nonrelativistic approach). The solution of the Faddeev equation leads to the following values of masses of the nonstrange baryons (S.Gerasyuta, A.Sarantsev, E.Yarevsky)

$m_N = 944$ MeV, $m_\Delta = 1222$ MeV. (15)

These masses as well as the calculated magnetic moments agree well with the experiment. However the calculated proton radius is noticeably lower than the experimental value

$R_p = 0.50$ fm, $R_p(\mathrm{exp}) \simeq 0.82$ fm. (16)

It should be stressed that all calculated proton radii in other potential models are essentially less than the experimental value.

5. LOW-LYING MESONS WITH HEAVY QUARKS (c,b)

Using bootstrap procedure we have calculated the structure of the quark-quark forces (fig. 12a). These forces can be use for obtaining the low-energy heavy quark interactions (see fig. 12 b,c). Here the structure (t-dependence) of the exchange states is the same but the vertex constants change.

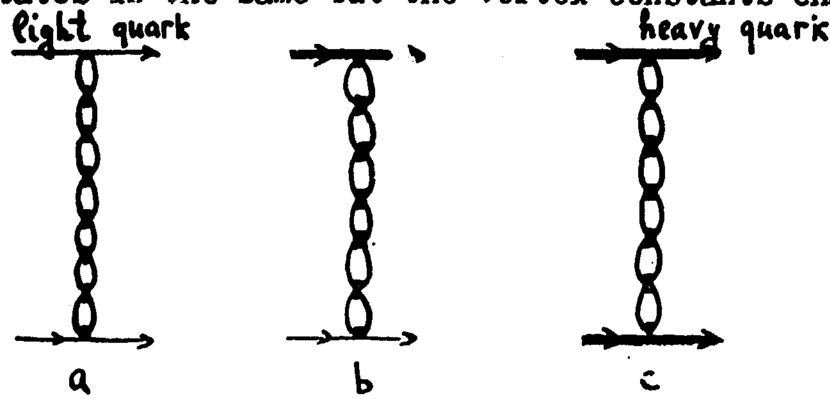

Fig. 12.

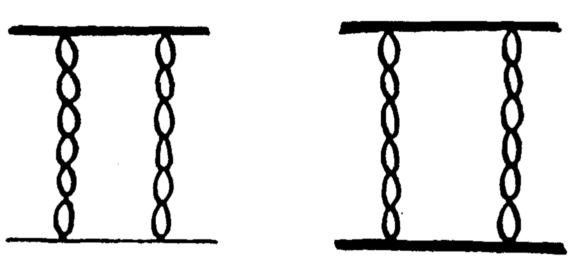

Fig. 13.

For the diagrams with multiple rescatterings of heavy quark (fig.13) we should change regularization parameters Λ as well.

The mass spectra of mesons containing c- and b-quarks were calculated using new vertices and cut-off parameters. The results are given in Table 3. The parameters(vertex constants and cut-off parameter) occurred to be smooth functions of the quark masses so they can be fitted by some simple formulae.

Using these extrapolation formulae one can calculate the masses of mesons with t-quark if the t-quark mass is fixed. For example, for m_t=80 GeV the binding energies $\mathcal{E}$ (where m(meson) = m_1(quark) + m_2(quark) - $\mathcal{E}$) are given in Table 4.

Table 3. (S.Gerasyuta, A.Sarantsev)

	0^- m (GeV)		1^- m(GeV)		0^+	
$u\bar{c}$, $d\bar{c}$	D (1.867)	1.867	D*(2.010)	2.010	-	2.119
$s\bar{c}$	D_s(1.971)	2.010	D_s^*(2.113)	2.120	-	2.300
$c\bar{c}$	η_c(2.980)	2.955	J/ψ(3.097)	3.097	χ_0(3.415)	3.453
$u\bar{d}$, $d\bar{b}$	B(5.270)	5.270	B*(5.320)	5.320	-	5.486
$s\bar{b}$	B_s(5.430)	5.375	B_s^*(5.390)	5.425	-	5.652
$c\bar{b}$	-	6.085	-	6.320	-	6.735
$b\bar{b}$	-	9.340	Υ(9.460)	9.460	-	10.071

Here the underlined masses are used for fixing parameters.

Table 4. (A.Sarantsev)

	0^- (GeV)	1^- (GeV)
$t\bar{t}$	7.23	7.15
$t\bar{b}$	0.08	0.01

6. CONFINEMENT FORCES

Up to now the processes with quarks were treated as if there is no confinement of the colour particles. Let us discuss qualitatively the character of the quark-quark forces which arise when confinement processes are taken into account.

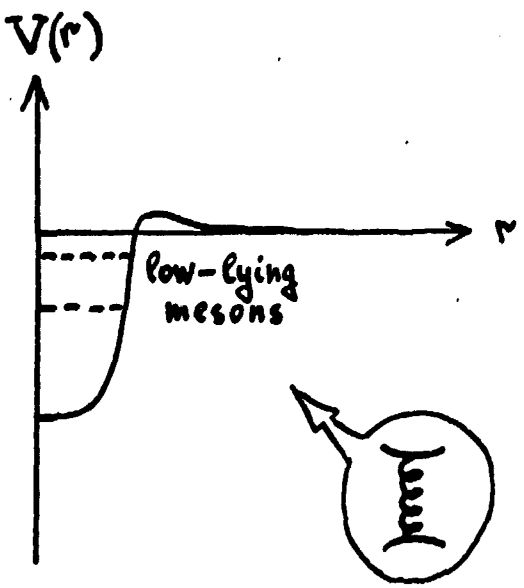

Fig. 14.

Bootstrap procedure has provided us with quark interaction where the constituent gluon played the main role. For visual demonstration let it be presented as potential V (fig. 14). This potential leads to the appearance of the low-lying mesons. But high-excited mesons are absent here: instead of discrete levels one has a continuous spectrum of $q\bar{q}$ states. This happens because processes which are typical for confinement are not taken into account. Namely, confinement of the $q\bar{q}$ pair with comparatively large energy is actually realized as the production of the new $q\bar{q}$ pair or several $q\bar{q}$ pairs (fig. 15).

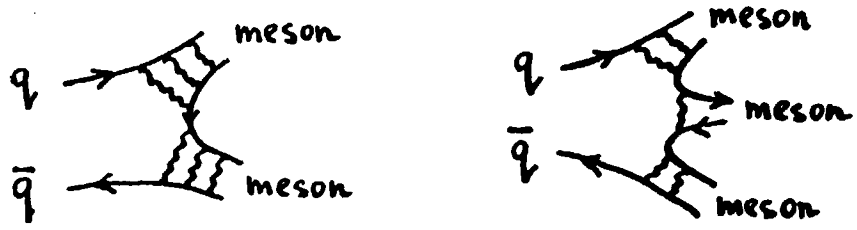

Fig. 15.

This means that in the transition $q\bar{q} \to q\bar{q}$ the forces appear which are connected with the processes of the fig. 16 type. Here in the t-channel the transitions of the octet colour states (c=8) play the main role. It means that these box-diagrams can

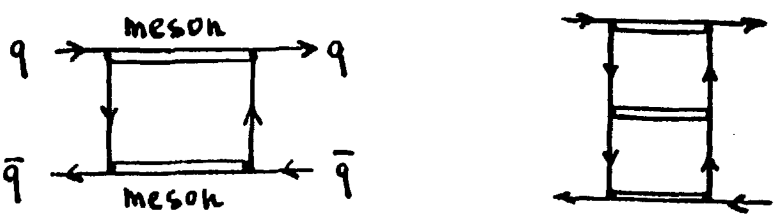

Fig. 16.

be important in the formation of hadron spectra. Let us draw qualitative picture for the potential of the $q\bar{q}$-interaction which is expected when confinement forces are considered(fig. 17). The potential barrier is determined by the box-diagrams.

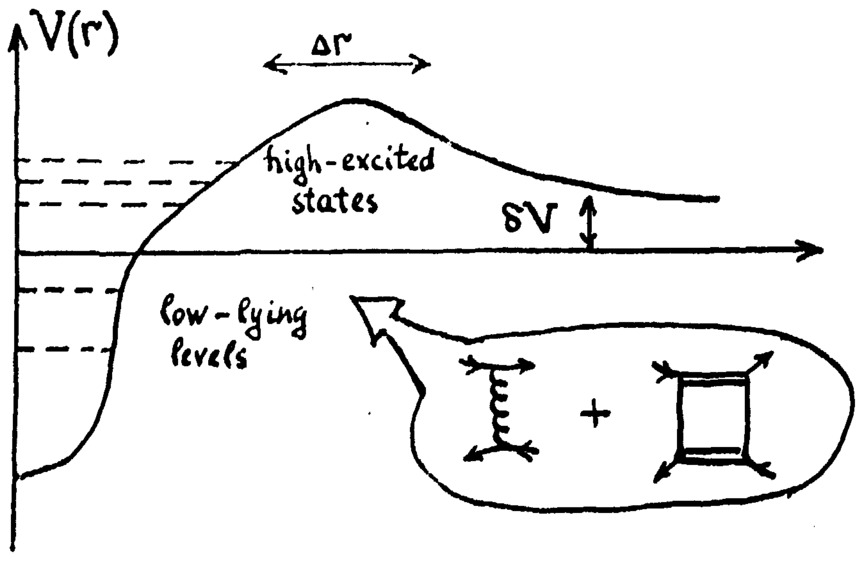

Fig. 17.

The barrier width (Δr) is given by the imaginary parts of these diagrams. The value of δV (or $V(r \to \infty)$) is determined by the lowest threshold

$$q_i q_j \to \text{meson } (q_i \bar{q}) + \text{meson } (q \bar{q}_j). \qquad (17)$$

Namely,

$$\delta V = m(q_i \bar{q}) + m(q \bar{q}_j) - m_i - m_j. \qquad (18)$$

Potential picture shown here is qualitative and rather rough:

(i) relativistic corrections for high-excited states are important (potential interaction is nonrelativistic one)

(ii) box-diagrams depend on energy (and lead in fact to 'quasipotential').

However we do not see any fundamental difficulty in taking into account the box-diagrams with the help of the dispersion technique which has been used here for description of low-energy states.

The picture discussed here is in qualitative agreement with the hypothesis of the supercritical nature of confinement forces made by Gribov [2].

We deal with the quarks as with real particles. However in soft region the quark diagrams should be treated as a spectral integrals over quark masses: the integration over quark masses in the amplitudes puts away the quark singularities and introduces the hadron ones.

7. SPECTRAL INTEGRATION OVER QUARK MASSES AS A REPRESENTATION OF THE DIAGRAMS WITH COLOUR PARTICLES IN THE REGION OF SOFT PROCESSES

Let us see an example of transition from Feynman diagram the quark-mass spectral representation. Consider the $q\bar{q}$ production by meson current (the ρ-meson current, for example - fig. 18a). If the confinement is neglected (i.e. confinement radius $R_c \to \infty$) one has for diagram of fig. 18a the following expression which depends on the quark mass m:

$$\Pi_{\alpha\beta}(p) = (-\delta_{\alpha\beta}\, p^2 + p_\alpha p_\beta)\, \Pi_m(p^2)$$

$$\Pi_m(p^2) = \int_{4m^2}^{\infty} \frac{dp'^2}{\pi} \frac{\Phi_m(p'^2)}{p'^2 - p^2 - i0} \qquad (19)$$

Here p is the momentum of the ρ-meson current and $\Phi_m(p^2)$ is imaginary part of $\Pi_m(p^2)$ determined by the cutting of the Feynman diagram of fig.18a. In the complex p^2-plane the self-energy

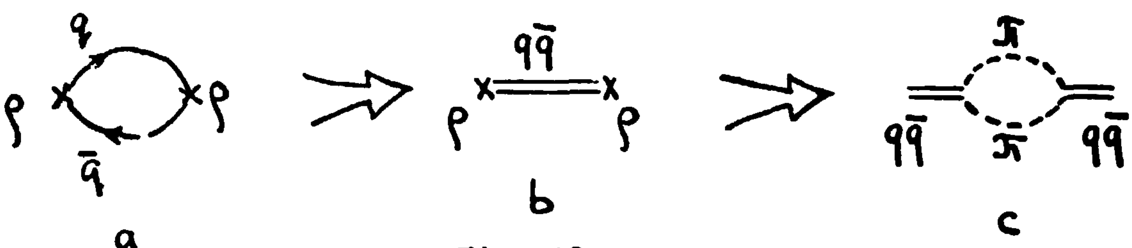

Fig. 18.

part $\Pi_m(p^2)$ has the threshold singularity $p^2 = 4m^2$ (fig. 19a). This threshold singularity is located on the first (physical) sheet of p^2-plane because we put $R_c \to \infty$.

Two phenomena arise when confinement radius R_c became finite. Firstly, the continuum of the $q\bar{q}$ states at $p^2 > 4m^2$ turns into a discrete set of states. It means that in the complex p^2--plane, instead of the cut from the point $p^2 = 4m^2$ (fig. 19a), one should have a set of poles (see figs. 18a, 19b). The discrete spectrum means that the dispersion integral (19) turns

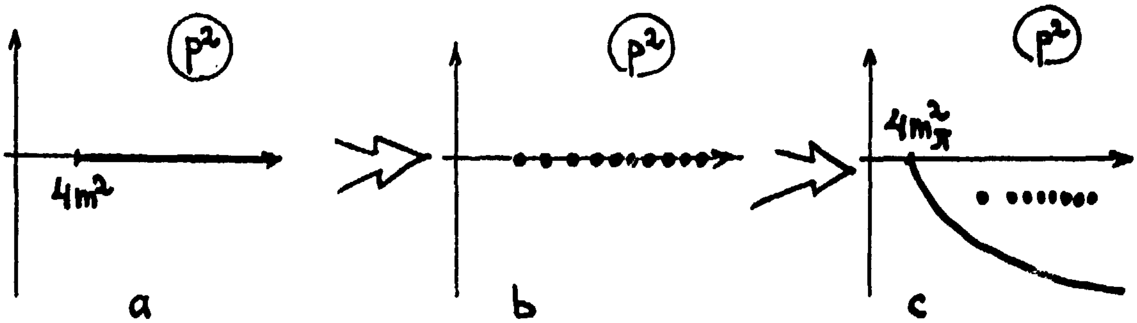

Fig. 19.

into the sum of the pole terms

$$\int_{4m^2}^{\infty} \frac{dp'^2}{\pi} \frac{\Phi_m(p'^2)}{p'^2 - p^2 - i0} \longrightarrow \sum_n \frac{\Phi_n}{p_n^2 - p^2} \tag{20}$$

Secondly, the energy-excited $q\bar{q}$-states produce the new $q\bar{q}$-pairs **due to** the confinement and decay into hadrons (fig. 18c **provides** the process $q\bar{q} \rightarrow 2\pi$ as an example).

The success of the quark model as well as the phenomenology **of the** soft hadron production argue in favour of the weak **dependence** of the quark amplitudes on the confinement. In this **case** it is possible to calculate the transition $q\bar{q} \rightarrow$ hadrons **using** perturbative method. Then in the first approximation we **have** for the mass-shift of the n-th level $\Delta p_n^2 = A_{nn}(p^2)$ **where** $A_{nn}(p^2)$ is the contribution of the graph of fig. 18c. **So we have**

$$\sum_n \frac{\Phi_n}{p_n^2 - p^2} \longrightarrow \sum_n \frac{\Phi_n}{p_n^2 - p^2 - A_{nn}(p^2)} \tag{21}$$

The $A_{nn}(p^2)$-amplitude has a threshold singularity at $s=4m_\pi^2$, **hence,** $\mathrm{Im}A_{nn}(p^2) > 0$ at $p^2 > 4m_\pi^2$. Therefore each pole of **the** series (21) is shifted onto the second sheet of the p^2-**-plane,** connected with the singular point $p^2 = 4m_\pi^2$ (fig,19c).

If the density of levels in eq. (21) is large enough it is **possible** to return to the integration over p'^2. **Assuming** also **that** $A_{nn}(p^2)$ depends slightly on n we have

$$\sum_n \frac{\Phi_n}{p_n^2 - p^2 - A_{nn}(p^2)} \longrightarrow \int_{p_{thr}^2}^{\infty} \frac{dp'^2}{\pi} \frac{\Phi_m(p'^2)}{p'^2 - p^2 - A(p^2)} \tag{22}$$

The amplitude $A(p^2)$ corresponds to the diagrams of the type **shown** in fig.16. The self-energy part $\Pi(p^2)$ satisfies the **confinement** condition: on the first (physical) sheet of the p^2-plane there are only hadron singularities whereas quark **singularities** are hidden in the second sheets (under the hadron cuts).

The eq. (22) gives the possibility to introduce the spectral integration over quark mass. Distribution density over quark mass, $\rho(m^2)$, is determined by the following way /3/

$$\Pi(p^2) = \int_{p^2_{thr}}^{\infty} \frac{dp'^2}{\pi} \frac{\Phi_M(p'^2)}{p'^2 - p^2 - A(p^2)} = \int_{m^2_\pi}^{\infty} dm^2 \rho(m^2)\, \Pi_m(p^2) \qquad (23)$$

Let us suppose that $\rho(m^2)$ is universal function. In this case more complicated quark diagrams, for example, quark triangle diagram of fig.20, is determined by the same quark density

$$T(q^2) = \int_{m^2_\pi}^{\infty} dm^2 \rho(m^2)\, T_m(q^2). \qquad (24)$$

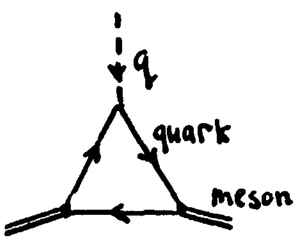

Fig. 20.

Here $T_m(q^2)$ is the amplitude for the process of fig.20 when the confinement is not taken into account ($R_c \to \infty$). The eq.(23) allows us to see directly that the amplitudes written as the spectral integrals over quark mass satisfy the requirement of the confinement. From the point of view of the analytical properties of the hadron amplitudes the absence of the quark confinement means the existence of the quark singularities corresponding to the production and the absorption of the real quarks in intermediate states. On the contrary, if there is quark confinement, hadron amplitudes should have on the physical sheets only the singularities which corresponds to the production and the absorption of real hadrons. Replacement of the quark singularities by the hadron ones can be carried out if the integration over quark masses is introduced. Integration over m eliminates quark singularities and introduces the hadron ones due to singularities of $\rho(m^2)$. For the light quarks $\rho(m^2)$ has a maximum near 300-400 MeV (the mass of the constituent quark).

We know that the rough approximation, when one neglects the quark mass distribution, i.e. when quark mass is fixed,

$$\rho(m^2) = \delta(m^2 - m_q{}^2) \tag{25}$$

is satisfying for low-lying hadrons. The problem, if the quark mass distribution is needed for high-excited states, requires a special consideration.

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