STATUS OF QCD

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A b s t r a c t

Status of Quantum Chromodynamics is reviewed. Nonperturbative and perturbative aspects of QCD are discussed. In considering nonperturbative QCD a special attention is paid to description of hadrons based on the operator expansion (QCD sum rule) method: calculation of meson and baryon masses, determination of partial widths, of formfactors at intermediate momentum transfers and of statical properties of hadrons in constand external fields. The main topics discus- in perturbative QCD are: coherence effects in the multiple production of gluons and quarks, heavy quark fragmentation functions, exclusive processes at large Q^2.

I. Introduction

For several years reviews on quantum chromodynamics started with the words "QCD is the best candidate for strong interaction theory". Now this time has gone, the choice has been done and QCD is no more a candidate but the true theory of strong interactions. This implies that the whole strong interaction physics: hadron masses, their decay widths, cross sections etc., are contained in a single line - QCD lagrangian:

$$L = i\sum_q{}' \bar{\varphi}_q^a (\gamma_\mu \nabla_\mu + i m_q)\varphi_q^a - \frac{1}{4} G_{\mu\nu}^n G_{\mu\nu}^n \quad (1)$$

where

$$\nabla_\mu = \partial_\mu + i g \frac{\lambda^n}{2} A_\mu^n$$

$$G_{\mu\nu}^n = \partial_\mu A_\nu^n - \partial_\nu A_\mu^n - g f^{nml} A_\mu^m A_\nu^n$$

$a = 1,2,3$, $n=1,2...8$ are colour indeces, A_μ^n - is the colour gluon field and summation is made over all quark flavours q=u,d, s,c,b... . Since nowadays we are aware of the QCD lagrangian, the focus of attention of strong interaction physics investigations is another than that five or ten years ago. If that time the most interesting were various checks of QCD: check of asymptotic freedom and of scaling violation, measurement of gluon spin etc., the main problems confronting us now are the QCD determination of hadron properties, elucidating the structure of the theory itself, in particular, of the vacuum structure, establishing the relationship between QCD and various model methods, and, of course, - the last but not least - the proof (constructive, desirable) of confinement.

QCD is the strong interaction theory: even at the highest accessible in the foreseen future momentum transfers the coupling constant $\alpha_s(Q^2) = g^2(Q^2)/4\pi$ is not very small, $\alpha_s(Q^2) \gtrsim 0.1$. At not very large momentum transfers, besides perturbative effects, there appear nonperturbative ones which exceed them. Calculation of high orders of perturbative theory is unreal and because of the series divergence it may be even unreasonable. Exact taking into account of nonperturbative effects is even more unreal. What can be expected is the development of some approximate methods. That is why neither at the present time nor in the future one cannot expect a high accuracy of QCD predictions, such as, say, in QED. As a rule, (with rare though very important exceptions) the accuracy of such predictions is at one or few tens of per cent. But in my opinion, this is not the reason to be grieved. As not only that experiment is brilliant which has a high accuracy but that in which new result has been obtained with a sufficient confidence, so in QCD reliability of theoretical results controllable inside the theory itself is important first of all. And this is an essential difference of QCD from model approaches where usually one has to go off the framework of a model to control realibility of the results.

In the last few years we have a considerable progress in finding various physical parameters of hadrons using the nonperturbative method grounded on the operator expansion and on the assumption of existence of nontrivial vacuum expectation values of quark and gluon fields in QCD. A great number of results has been obtained on this way: masses, widths, formfactors at small and intermediate momentum transfers, hadron properties in constant external fields etc. were determined. The same method enables one to check various models. In the first part of my talk dedicated to nonperturbative QCD I review basic ideas and possibilities of such an approach and give the results obtained in this way.

In the second part I will dwell on new perturbative QCD results which are in my opinion the most important. The status of QCD will be discussed in my report from the standpoint of how QCD can describe observed physical phenomena and which conclusions upon the structure of QCD itself follow from this comparing with experiment. In this connection I will not discuss a number of principal but now pure theoretical problems of QCD like confinement, phase transitions, relationship between QCD and string theory etc.

II. Nonperturbative QCD.

I. Light Hadron Masses.

The method in view which is often called the sum rule QCD method has been originally suggested by Shifman, Vainshtein and Zakharov [1] and has been applied by them for determining masses and leptonic widths of light mesons(ρ , π , A_1,K^*) and for studying some parameters of charmonium. I dwell here on considering the properties of hadrons comprised of light quarks where the most essential progress has been achieved in the last few years.

The method is grounded on the following considerations:

(i) in the virtuality region of order $Q^2 \sim 1$ GeV^2 the strong interaction constant $\alpha_s(Q^2)$ is already rather small, $\alpha_s(Q^2)\sim 0.3-0.4$, so that nonperturbative terms are small $\alpha_s/\pi \sim 0.1$ and the leading logarithmic corrections $\sim [\alpha_s(Q^2)\ln Q^2/\Lambda^2]^n$ can be easily taken into account.

(ii) the nonperturbative effects which reduce to appearance of vacuum condensates play a fundamental role. Of the vacuum condensates the most essential are the quark condensate density $\langle 0|\bar{q}q|0\rangle$, q=u,d,s, and the gluon condensate density $\langle 0|G^n_{\mu\nu}G^n_{\mu\nu}|0\rangle$. Vacuum condensates in this approach are considered as phenomenological parameters determined either from experiment or from selfconsistency of the sum rules obtained.

It is convenient to present characteristic features of the method on an example of proton mass calculation [2,3] to which I now turn. Consider the polarization operator

$$\Pi(p)=i\int d^4x\, e^{ipx}\langle 0|T\{\eta(x),\bar{\eta}(0)\}|0\rangle \quad (2)$$

where $\eta(x)$ is the quark current with proton quantum numbers and p^2 is chosen to be space-like: $p^2 < 0$, $|p^2| \sim 1$ GeV^2 . Current η is the colourless product of three quark fields $\eta = q^a q^b q^c \varepsilon^{abc}$, q=u,d, the form of the current will be specialized below. The general tensor structure of Π is

$$\Pi(p)=\hat{p}f_1(p^2)+f_2(p^2) \quad (3)$$

For each of the functions $f_i(p^2)$, i=1,2 the following operator expansion can be written:

$$f_i(p^2)=\sum_n{}' C_n^{(i)}(p^2)\langle 0|O_n^{(i)}|0\rangle \quad (4)$$

where $\langle 0|O_n^{(i)}|0\rangle$ are vacuum expectation values (v.e.v) of different operators (vacuum condensates), $C_n^{(i)}(p^2)$ are calculated in QCD functions.

As is known, the current u- and d-quark masses entering the lagrangian of QCD (1) are very small - of order several MeV, so they can be neglected with a very good accuracy, i.e. u- and d-quarks can be taken massless. Then accounting for only u- and d-quarks, the QCD lagrangian is chirally-invariant and if chiral invariance would not be spontaneously broken, function $f_2(p^2)$ would be identically zero. In reality, however, the chiral invariance is spontaneously broken in QCD. The first evidence of this is, of course, the existence of baryon masses. Another signal of chiral symmetry breaking is the appearance of chiral violating v.e.v's, the most important of which is the quark condensate density. The question arises, how these two phenomena are connected and, particularly, if the proton mass can be expressed through $\langle 0|\bar{q}q|0\rangle$. The answer to the last question is affirmative and a bit later I shall demonstrate the corresponding formula. Since $\langle 0|\bar{q}q|0\rangle$ is the lowest in dimension chirality violating operator, the operator expansion for $f_2(p^2)$ starts from the term proportional to $\langle 0|\bar{q}q|0\rangle$.

To demonstrate how the operator expansion (4) is built I present several first terms of it classifying the terms according to the operator dimension. The zero term of of the operator expansion with d=0 corresponding to the unit operator is described by the graph of fig.1a and has the form

$$C_0 \hat{p} p^4 \ln \frac{\Lambda_u^2}{-p^2} + polynom, \qquad (5)$$

where C_0 is a constant, Λ_u is ultraviolet cut-off.

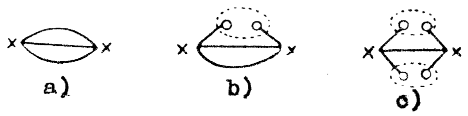

Fig.1.

The main diagrams accounted for calculation of $\Pi(p)$ (2). Solid lines correspond to quarks, circles outlined by dots stand for v.e.v.'s of field operators.

This term preserves chiral invariance and contributes to function $f_1(p^2)$. The chirality violating operator qq with d=3 is the next in dimension. Its contribution to $f_2(p^2)$ is described by the graph of fig. 1b and is equal to

$$C_1 \langle 0|\bar{q}q|0\rangle \ln \frac{\Lambda_u^2}{-p^2} + polynom \qquad (6)$$

The most essential correction in the operator expansion for $f_1(p^2)$ arises from the four-quark operators $\bar{q}\Gamma_i^{(K)}q\cdot\bar{q}\Gamma_i^{(K)}q$ with d=6 whose contribution is described by the graph of fig.1c and has the form

$$\sum_K{}' C_2^{(K)} \langle 0|\bar{q}\Gamma_i^{(K)}q\cdot\bar{q}\Gamma_i^{(K)}q|0\rangle \frac{\hat{p}}{p^2} \qquad (7)$$

This contribution is very important numerically since unlike the graph of fig.1a containing two-loop integration, there is no integration in the graph of fig.1c, so that eq.(7) contains, compared to (5), a large numerical factor: $C_2/C_0 \sim (2\pi)^4$. When calculating v.e.v (7) there is often used the factorization hypothesis according to which in expansion of the four-quark operator product over intermediate states the main contribution is given by the vacuum state. Arguments based on the $1/N_c$ expansion are in the favour of this hypothesis. The above three terms of the operator expansion indicate the way of calculating $\Pi(p)$ in QCD, i.e. the way of finding the left-hand part of the desired sum rules. On the other hand, functions $f_i(s)$, $s=-p^2$ may be expressed via the characteristics of physical states using dispersion relations

$$f_i(s) = \frac{1}{\pi}\int_0^\infty \frac{Im f_i(\rho^2)}{\rho^2+s}\, d\rho^2 + polynom \qquad (8)$$

It is useless to equate (4) and (8) since both the left-hand and right-hand parts of resultant equality contain unknown polynomials. In order for the appeared equality to acquire the meaning one has to apply to both its sides the Borel (Laplace) transformation [1] defined as

$$B_{M^2} f(s) = \lim_{s\to\infty,\ n\to\infty,\ s/n=M^2=Const} \frac{s^{n+1}}{n!}\left(-\frac{d}{ds}\right)^n f(s) = \frac{1}{\pi}\int_0^\infty e^{-\rho^2/M^2}\, Im f(\rho^2)\, d\rho^2 \qquad (9)$$

if $f(s)$ is given by dispersion relation (8). Notice that

$$B_{M^2}\frac{1}{s^n} = \frac{1}{(M^2)^{n-1}(n-1)!} \qquad (10)$$

The Borel transformation permits to attain three goals at once:

(1) to nulify subtraction terms;

(2) to suppress the contribution of the highest excited states compared to the contribution of the requisite lowest state (proton);

(3) to suppress contributions of high order terms in the operator expansion (owing to factor $1/(n-1)!$ in (10)).

The lowest state (proton) contribution into imaginary part of $\Pi(p)$ has the form

$$Im\,\Pi(p)_p = \pi\langle 0|\eta|p\rangle\langle p|\bar{\eta}|0\rangle\,\delta(p^2-m^2) = \lambda_N^2\pi(\hat{p}+m)\,\delta(p^2-m^2) \qquad (11)$$

where

$$\langle 0|\eta|p\rangle = \lambda_N v(p)$$

λ_N is a constant and v is the proton spinor. It is clear from (11) that the proton contribution will dominate in some region of the Borel parameter M^2 only in the case when both QCD calculated functions $mf_1(M^2)$ and $f_2(M^2)$ are of the same order. This means that the value of quark condensate must explain the numerical value of proton mass.

To improve and control the accuracy in dispersion representation (8) one should also take into account the highest state contribution. It is usually made replacing $\mathrm{Im}\, f(p^2)$ by contributions of the simplest quark loops (fig.1) starting from some "continuum threshold" W.

It should be emphasized that in the sum rule method the presence of structure in hadronic spectra in the small mass region, i.e. the appearance of resonances separated by a dip from the region of smooth continuum which corresponds to parton model, is not inserted into theory outside but follows from existence of power corrections.

A few words on the choice of the quark current $\eta(x)$. In case of baryons (unlike mesons) even if we restrict ourselves by currents without derivatives, there exists as a rule several currents with quantum numbers of a given baryon. The choice between them should be done from physical reasons in order to provide: (1) renormcovariance; (2) existence of nonrelativistic limit; (3) the above formulated requirement (for proton) for the functions f_1 and f_2 to be of the same order; (4) covergence of the operator expansion series within accounted terms. In case of proton all these requirements are satisfied by the current [2, 4] (see also [5])

$$\eta = (u^a C\gamma_\mu u^b)\gamma_\mu \gamma_5 d^c \cdot \varepsilon^{abc} \qquad (12)$$

The mentioned above generally describes how the sum rule for proton mass determination is found. In the calculation of ref. [6] the operators up to dimensions d=9 were taken into account.

Owing to the factorization hypothesis two sum rules corresponding to functions $f_1(p^2)$ and $f_2(p^2)$ are expressed through three v.e.v's: $\langle 0|\bar q q|0\rangle$, $\langle 0|(\alpha_s/\pi)\, G^n_{\mu\nu} G^n_{\mu\nu}|0\rangle$, $-g\langle 0|\bar q\, \sigma_{\mu\nu}(\lambda^n/2)\, G^n_{\mu\nu}\, q|0\rangle$.

The proton mass m and the proton transition constant into the quark current λ_N (W is also a variable parameter) may be found from the sum rules by the best fitting. Such a fitting should be made within a restricted interval of M^2 where, on one hand, the continuum contribution is rather small (say, less than 50%) which restricts M^2 from above and, on the other hand, highest power corrections are small (say, $< 10\%$) which restricts M^2 from below. Off this interval the accuracy of the theory is incontrollable, for instance, at very large M^2 the lowest state (proton) contribution is hidden in the background and the result for m and λ_N depends cardinally on continuum model. I wish to emphasize that without estimating the contribution of highest states and highest power corrections the results obtained by the QCD sum rule method cannot be considered as reliable. The best fitting in the permissible interval $0.7 < M^2 < 1.2$ GeV^2 at the chosen values of v.e.v's (for discussion of their numerical values see below) and W=1.5 GeV gives [2,6] for the proton mass m = 1.0±0.1 GeV. For m there is also a simple approximate formula

$$m = \left[-(2\pi)^2 \cdot 2\langle 0|\bar q q|0\rangle\right]^{1/3} \qquad (13)$$

which reflects the fact the appearance of proton mass is connected with spontaneous violation of chiral invariance - i.e. with the presence of quark condensate.

The accuracy of the obtained sum rules can be checked by substituting into it the experimental value of m and plotting the M^2 dependence of λ_N^2 . The results are presented in fig.2 where $\tilde\lambda_N^2 = 32\pi^4\lambda_N^2$, solid line corresponds to sum rule for $f_1(p^2)$, dashed line - for $f_2(p^2)$.

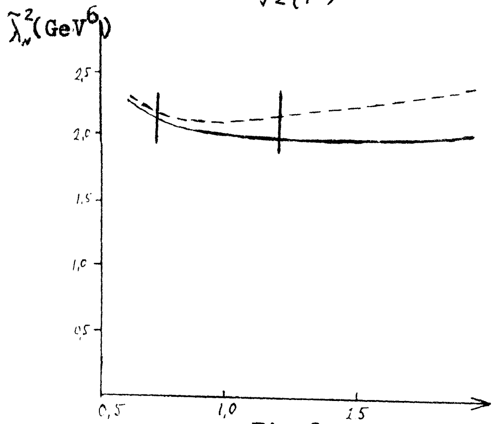

Fig.2

As is seen from fig.2 the difference of the two determined sum rules and deviation of $\tilde\lambda_N^2$ from the constant within $0.7 < M^2 < 1.2$ GeV^2 are less than 10%. Such an accuracy is natural because highest power corrections are in this case of the same order. As a result we have (at the normalization point μ = 0.5 GeV)

$$\tilde\lambda_N^2 = 32\pi^4\lambda_N^2 = 2.1 \pm 0.2 \text{ GeV}^6 \qquad (14)$$

I wish to remark that in the case of baryons (unlike the meson case) there are several sum rules (for baryons with spin 1/2 two, for baryons with spin 3/2 - seven) which all must give the same values of the baryon mass and residue (in the limit of accuracy of each sum rule). This is an important check of the whole approach.

Masses and leptonic widths of meson decays and masses of other baryons are calculated analogously. Not dwelling on details, I give the calculation results (Tables I and 2).

Table I.

Meson (ref.)	J^{PC}	Mass (GeV) theor.	Mass (GeV) exp.	Lept. decay const. theor.	Lept. decay const. exp.
ρ [1]	1^{--}	0.78 ± 0.04	0.770	$g_\rho^2/4\pi = 2.4 \pm 0.2$	2.54 ± 0.23
ω [1]	1^{--}	0.78 ± 0.04	0.783	22 ± 2	18.4 ± 1.8
φ [1]	1^{--}	1.07 ± 0.05	1.02	14	11.7 ± 0.9
K^* [1]	1^{--}	0.93 ± 0.05	0.892	1.4	-
π [2]	0^{-+}	-	0.140	$f_\pi = 125$ MeV	133 MeV
A_1 [7]	1^{++}	$1.25 \pm .15$*)	1.27.	$g^2/4\pi \approx 6$	≈ 6
D [7]	1^{++}	1.25 ± 0.15*)	1.28	$g^2/4\pi \approx 6$	
f [8]	2^{++}	1.25 ± 0.05	1.27	$g_f = 0.040$; $\Gamma_{f\pi\pi} = 200 \pm 30$ MeV	-; 180 MeV
A_2 [8]	2^{++}	1.25 ± 0.05	1.32	-	-
A_3 [8]	2^{-+}	1.63 ± 0.1	1.68	-	-

*) Corrected by the speaker

Predictions for η and η' and for scalar mesons are absent in Table I: QCD sum rules do not work in the case of scalar and pseudoscalar channels as well as in the case of longitudinal axial channel with isospin zero. In particular, the sum rule approach cannot explain the 100% violation of the Okubo-Zweig-Iizuki rule in the pseudoscalar (or in the longitudinal axial) channel with T=0, i.e. the fact that η -meson is an octet component and η' a unitary singlet. A possible (though not proved) explanation of this fact is in an important effect of direct instantons.[9, 10] Though scalar and pseudoscalar channels cannot be quantitatively considered by the sum rule method one can get an important quantitative conclusion [10] : the characteristic mass scale (i.e. the distance between resonances) in these channels is much more large (by a factor of 3 in m^2) than in the vector channel.

Table 2.

Baryons

Baryon	J^P	T	Calculated value (GeV)	Theor.[2,6]	Exp.
N	$1/2^+$	1/2	m_N	1.0 ± 0.1	0.94
Δ	$3/2^+$	3/2	m_Δ	1.37 ± 0.15	1.23
N^*	$3/2^-$	1/2	m_{N^*}	1.75 ± 0.25	1.52
Λ	$1/2^+$	0	$m_\Lambda - m_N$	0.19	0.175
Σ	$1/2^+$	I	$m_\Sigma - m_N$	0.23	0.25
Ξ	$1/2^+$	1/2	$m_\Xi - m_N$	0.40	0.38
Σ^*	3/2	I	$m_{\Sigma^*} - m_\Delta$	0.14	0.15
Λ^{**}	$3/2^-$	0	$m_{\Lambda^{**}} - m_{N^*}$	0.14	0.17
Σ^{**}	$3/2^-$	I	$m_{\Sigma^{**}} - m_{N^*}$	0.18	0.15
Ξ^{**}	$3/2^-$	1/2	$m_{\Xi^{**}} - m_{N^*}$	0.30	0.30

Remark. The mass differences are calculated in the linear approximation in $m_Y - m_N$, m_s, and f. The values m_s and f are chosen in order to obtain a best fit of experimental data (m_s=0.105 GeV, f =-0.11).

When calculating meson and baryon masses without strange quarks, only the above mentioned v.e.v.'s enter. In considering hyperons and mesons with strange quarks there appear two new parameters - the strange quark mass m_s and the parameter $f = \langle 0|\bar{s}s|0\rangle / \langle 0|\bar{u}u|0\rangle - 1$. In order to find m_s and f with a higher accuracy one should substitute into sum rules the experimental values of hyperon masses and find m_s and f from the best fit (not exploiting the expansion in mass differences $m_Y - m_N$). This gives (at μ =0.5 GeV)

$$m_s = 150 \pm 30 \text{ MeV} \quad f = -0.2^{-0.05}_{+0.1} \tag{15}$$

Now about the parameter λ_N which determines the proton transition amplitude into quark current η . As was seen from (14) and from fig.2 λ_N can be found from sum rules with a good accuracy. Exact knowledge of the constant λ_N and of other constants of the same type is of great importance since they: (1) are used in calculating baryon magnetic moments (see below); (2) determine the proton life time in the grand unification SU(5) theory; (3) determine the constant in the asymptotics of the neutron electromagnetic formfactor. The latter point needs a more detailed explanation. The neutron electromagnetic formfactor in the asymptotics at $Q^2 \to \infty$ has the form [11, 12]

$$F_{1n}(Q^2)_{as} = \left[4\pi\alpha_s(Q^2)\right]^2 \left[\frac{\alpha_s(Q^2)}{\alpha_s(\mu^2)}\right]^{\frac{4}{27}} \frac{100}{3} \frac{h^2}{Q^4} \quad (16)$$

where h is determined by proportional to p_μ part of the matrix element

$$\varepsilon^{abc}\langle 0|(u^a C\gamma_\mu u^b)d^c|p\rangle = (p_\mu h + \gamma_\mu \beta)v \quad (17)$$

The matrix element (17) was determined from the sum rules [6] and it was found that h = = 0,8.10^{-2}GeV^2 . Substitution of h into (16) even at the largest experimentally assecсible Q^2=25 GeV^2 gives $F_{1n}(Q^2)_{Q^2=25\,GeV^2}$ = 1,2.$10^{-2}/Q^4$, a value by two orders of magnitude smaller than the experimental one, $F_{1n}^{exp} \approx -1/Q^4$ (and of the opposite sign). Hence it follows that asymptotic formulae for the nucleon electromagnetic formfactor have nothing to do with that is being observed (or will be observed) at accessible Q^2 . The analogous conclusion has been obtained by Isgur and Llewellyn-Smith [13] on the basis of the nucleon quark wave function model. Constants λ_N can be also used for checking and correcting the bag model. Since in the bag model $\lambda_N^2 \sim |\psi(0)|^2 \sim R^{-6}$ where R is the bag radius, the knowledge of λ_N^2 enables one to find R with a high accuracy if other parameters are fixed [14] .

The situation with glueballs is still far from being clear since in QCD only restricted statements can be done in a modelless way. In the sum rule method for the case of scalar and pseudoscalar glueballs, the direct instanton contribution seems to exist which makes unreliable the results obtained with this method. If it is the case, one should expect, owing to the same mechanism, a strong mixing of 0^+ and 0^- gleuballs with corresponding quark states. There are some arguments [10] that the 0^+ glueball mass is rather small $m_g(0^+) \lesssim 1.4$ GeV, while the 0^- glueball must be noticeably heavier, $m_g(0^-) \sim 2$-2.5 GeV. For the tensor 2^{++} glueball the statements are more definite: its mass is within 1.5-2.0 GeV [10, 15] and since direct instantons do not contribute to the corresponding sum rules the mixing of 2^{++} glueball with quark states should not be large especially. In this case, however, as is shown by the calculations [15] , the continuum threshold is close to the lowest state mass, i.e. one cannot expect a resonance well separated from the background.

A few words on the hybrid quark-antiquark gluonic mesons $\bar{q}qG$ (meiktons). QCD sum rules predict [16, 17] for the hybrid meson mass with spin 1 to be equal to m=1.2±0.2GeV. The main uncertainties in these calculations are due to a rather large continuum contribution and to essential role of purely known v.e.v's of high dimension operators. (Besides, there are discrepances in the results of refs. [16, 17]). Within these uncertainties it is difficult to say what is heavier, 1^+ or 1^- meson (though the authors of ref. [17] claim that 1^+ is lighter). Hybrid spin 0 mesons appear to be much more heavy, ~3 GeV. I will not discuss here the main problem of glueball and hybrid meson physics: how to find an appropriate signature for their observation (see, for example, review [18]).

2. Formfactors and Hadronic Widths.

The sum rule method was generalized to consideration of hadronic formfactors and widths [19, 20] . The initial point of such generalizing is studying of the vertex function

$$\Gamma_{AB}(p,p';q) = -\int d^4x\, d^4y\, \exp[i(p'x - qy)] \times \langle 0|T\{j_B^+(x), j(y), j_A(0)\}|0\rangle, \quad q = p' - p \quad (18)$$

where j_A, j_B are quark currents with quantum numbers of hadrons A and B and $j = j^{el}$ if we are interested in the electromagnetic formfactor $F_{AB}(q^2)$. $\Gamma_{AB}(p,p';q)$ is studied in the kinematic region $p^2 < 0$, $p'^2 < 0$, $q^2 = -Q^2 < 0$, at $|p^2| \sim |p'^2| \sim 1\,GeV^2$, $Q^2 \gtrsim 1$ GeV^2 . After this, on one hand, Γ_{AB} is calculated in QCD using the operator expansion and on the other hand, is represented via physical state contributions with the help of the double dispersion relation in p^2 and p'^2 . In QCD calculations two condensates $\langle 0|G_{\mu\nu}^n G_{\mu\nu}^n|0\rangle$ and $\alpha_s\langle 0|\bar{q}q|0\rangle^2$ are taken into account. The effect of power corrections and of excited state transition increases with Q^2 rising. The latter is evident since, as is well known, in deep inelastic scattering the relative role of inelastic transitions increases with increasing

Q^2. This circumstance restricts from above the Q^2 region where the approach in view is applicable. At small Q^2 this approach is inapplicable because at $q^2 \to 0$ in the t-channel large distances are important. The signal to it is the appearance of nonphysical singularities $\sim 1/Q^2$, $\ln Q^2$ etc. As a result the whole consideration appears to be valid in the region of intermediate Q^2, $0.5 \lesssim Q^2 \lesssim 3$ GeV^2. (One manage to obviate the restrictions at small Q^2 by introducing into play new v.e.v's - see below). Several hadronic formfactors were calculated on this way and I give here the results.

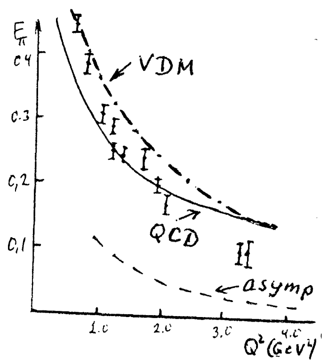

Fig.3.

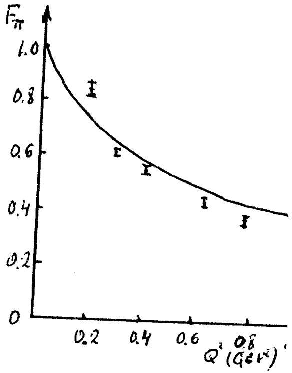
Fig.4.

Fig.3 represents the QCD found values of the pion formfactor [19] (solid curve) in comparison with experimental data. Fig.3 shows also the vector dominance model prediction (dot-dash curve) and the asymptotic QCD formula [21-25] (dashed curve).

$$F_\pi(Q^2) = \frac{8\pi\alpha_s(Q^2)}{Q^2} f_\pi^2 \quad , \quad f_\pi = 133 \text{ MeV} \tag{19}$$

The QCD calculated curve agrees with experiment in the region where QCD formulae are valid. The VDM curve differs a little from QCD and also sufficiently describes the experiment. However, the asymptotic formula (19) in the region $Q^2 \sim 3$ GeV^2 results in the values of $F_\pi(Q^2)$ by few times smaller than QCD (and experiment) gives. Fig.4 shows the results of the recent calculations [26] of the pion electromagnetic formfactor at small $Q^2 < 1$ GeV^2. In these calculations new quantities appear: values of vacuum condensates in the presence of external electromagnetic field. I will discuss the quantities of this type in the next Section. The calculations of ref. [19] and [26, 108] (the first one by analytical continuation $F_\pi(Q^2)$ to the point $Q^2=0$ using dispersion relation in Q^2 and information on hadronic spectrum and asymptotics) give precisely the same value of the pion radius $r_\pi^2 = 0.44 \pm 0.04 f^2$ in an excellent agreement with the recent data [27] $r_\pi^2 = 0.43 \pm 0.04.\ fm^2$

Fig.5 exhibits the QCD calculated [19] formfactors of ρ-meson : electric $F_1(Q^2)$, anomalous magnetic $F_2(Q^2)$ and quadrupole $F_3(Q^2)$ as well as the ratio $F_2(Q^2)/F_1(Q^2)$ (dashed curve, the right-side scale). Of attention in Fig.5 is that the ratio $F_2(Q^2)/F_1(Q^2)$ can be rather well extrapolated to the point $Q^2=0$ and $F_2(0)/F_1(0) \approx 1$. The property $F_2(0)/F_1(0)=1$, i.e. that the gyromagnetic ratio for ρ- meson is 2, is a characteristic feature of VDM in which ρ-meson is considered as the Yang-Mills boson [28-30]. Thereby at small Q^2 QCD surprisingly confirmes VDM at this nontrivial point.

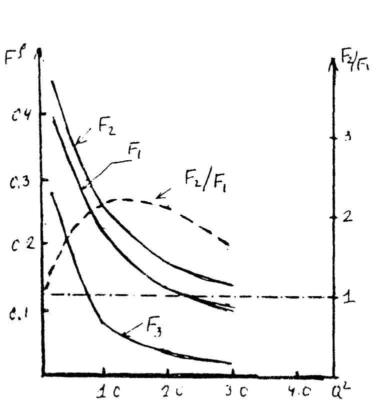

Fig.5.

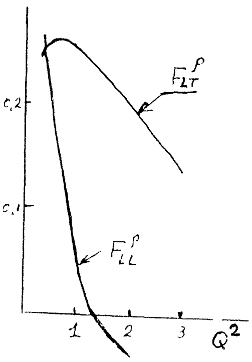

Fig.6.

One can also compare the ρ-meson formfactor behaviour with expectations of asymptotic QCD. To this end it is convenient to plot graphs for longitude-longitudinal $F_{LL}(Q^2)$ and longitudinal-transverse $F_{LT}(Q^2)$ formfactors (fig.6). It is expected in the asymptotics [21, 31] that $F_{LL}(Q^2) \sim 1/Q^2$, $F_{LT}(Q^2) \sim 1/Q^3$, i.e. $F_{LL} >> F_{LT}$. We see that at intermediate Q^2 the inverse relation takes place.

In ref. [32] the electromagnetic $\rho \to \pi$ transition formfactor in the intermediate region $0.5 < Q^2 < 3$ GeV^2 has been calculated. Now and again, there is a good coincidence between QCD and VDM results

$$F_{\rho\pi}(Q^2)_{VDM} = (g_{\omega\rho\pi}/g_\omega)\left[m_\omega^2/(Q^2+m_\omega^2)\right]$$

and strong disagreement with the asymptotic QCD where $F_{\rho\pi}(Q^2) \sim 1/Q^4$ is expected.

Knowledge of formfactors in the region of intermediate Q^2 enables one, using dispersion relation for $F(Q^2)$ in Q^2 and saturating it by the lowest states contributions, to determine some interaction constants and, respectively, hadronic widths. The following parameters were found in this way:

$$\frac{g_\rho^2}{4\pi} = 3.4\pm0.3 \quad [19] \quad \Gamma_{\rho\pi\pi} = 175\ \text{MeV} \pm 10\% \quad (\Gamma^{exp}_{\rho\pi\pi} = 155\pm5\ \text{MeV}) \tag{20}$$

$$g_{\omega\rho\pi} = 17\pm3\ \text{GeV}^{-1} \quad [33] \quad \Gamma_{\omega3\pi} = 11.4\ \text{MeV} \pm 30\% \quad (\Gamma^{exp}_{\omega3\pi} = 8.9\ \text{MeV}) \tag{21}$$

The same technique was used [34] to calculate the formfactor of transitions $D^+ \to \bar{K}^0 e^+\nu$ and to determine its value at $q^2=0$. It was found that $f_+^{D\to K}(0)=0.6\pm0.1$ what is definitely less than usually accepted value $f_+^{D\to K}(0)=1$. The knowledge of $f_+^{D\to K}(0)$ by comparing with experimental data on the $D^+\to\bar{K}^0 e^+\nu$ width permits one to improve the lower limit on element $|U_{cs}|$ of the Kobayashi-Maskawa matrix.

Finally, I will demonstrate the curves of proton and neutron magnetic formfactors in the region of small $Q^2 < 1$ GeV^2 obtained in the recent paper [35] (figs.7,8).

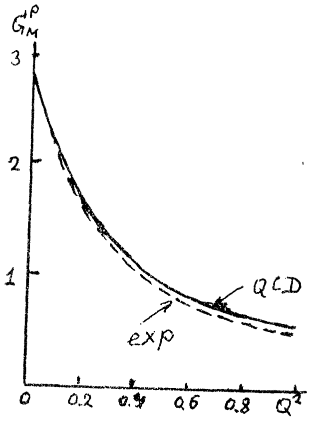

Fig.7.

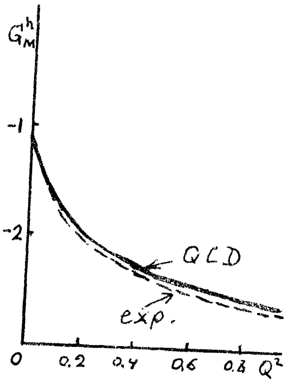

Fig.8

Formfactors G_M^p and G_M^n are calculated employing a rather different technique than discussed above: to calculate them one should know field induces v.e.v's, particularly, the discussed below quantity $\chi(q^2)$. The advantage of the method used in [35] is that it makes it possible to pass into the small Q^2 region. As is seen from figs. 7,8, the coincidence of the calculations with experiment is good enough.

3. Hadrons in External Fields.

The method presented in the preceding Section is inapplicable to consideration of hadron characteristics in static (or slow varying) external fields, for instance, to calculation of nucleon magnetic moments. The physical reason of this is clear: one must take into account the influence of external field on the properties of vacuum. As a result of such effect there appear new, previously absent, field induced v.e.v's. E.g., $\langle 0|\bar{q}\sigma_{\mu\nu}q|0\rangle = 0$ in the absence of external fields. However, if quarks are in the weak constant external electromagnetic field $F_{\mu\nu}$, then in general case it can be written

$$\langle 0|\bar{q}\sigma_{\mu\nu}q|0\rangle_F = \sqrt{4\pi\alpha}\,\chi_q F_{\mu\nu}\langle 0|\bar{q}q|0\rangle \tag{22}$$

The coefficient χ_q has the meaning of magnetic susceptibility of quark condensate. (The factor $\langle 0|\bar{q}q|0\rangle$ is introduced into the r.h.s of (22) for a reason that $\langle 0|\bar{q}\sigma_{\mu\nu}q|0\rangle$ just as $\langle 0|\bar{q}q|0\rangle$ breaks chiral invariance). After introducing field induced v.e.v's (similar to (22)) the problem of calculating nucleon magnetic moments can be solved [36, 37] considering polarization operator of quark currents with nucleon quantum numbers (2) and assuming that quarks move in the constant electromagnetic field $F_{\mu\nu}$. The only prise which has to be paid is the assumption that χ_q is proportional to the quark charge e_q, $\chi_q = e_q\chi$, i.e. that the contribution of graph 9b is small compared to the graphs of fig.9a. (Some arguments can be given why it is really so [36]).

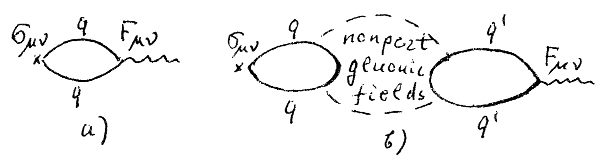

Fig.9.

When calculating nucleon magnetic moments by the QCD sum rule method (as well as in any other problems in constant external fields) there appears a difficulty stemming from the fact that in describing polarization operators in terms of physical states, besides the interesting for us one-nucleon contribution proportional to $\lambda_N^2\mu/(p^2-m^2)^2$ there exist contributions which correspond to transitions into excited states $N\to N^*$ proportional to $1/(p^2-m^2)(p^2-m^{*2})$. After Borel transformation the contributions of both types have one and the same exponent $exp(-m^2/M^2)$, i.e. $N \to N^*$ transitions do not turn to be exponentially suppressed. They must be taken into account exactly but one manage to get rid of them by studying the M^2 dependence of sum rules.

I will give here the sum rules corresponding to structures $\hat{p}\sigma_{\mu\nu} + \sigma_{\mu\nu}\hat{p}$ and $i(p_\mu\gamma_\nu - p_\nu\gamma_\mu)\hat{p}$ found in [36] for calculating proton and neutron magnetic moments

$$\mu_p e_d - \mu_n e_u = \frac{4a^2}{3\tilde{\lambda}_N^2}(e_u^2 - e_d^2)\left[1 - M^2\frac{\partial}{\partial M^2}\right]e^{\frac{m^2}{M^2}} L^{4/9} \quad (23)$$

$$\mu_p e_u - \mu_n e_d = e_u + \frac{4am}{\tilde{\lambda}_N^2}(e_u^2 - e_d^2)\left[1 - M^2\frac{\partial}{\partial M^2}\right] M^2 exp(m^2/M^2) \quad (24)$$

The notations are $a = -(2\pi)^2\langle 0|\bar{q}q|0\rangle$, $L = \ln(M/\Lambda)/\ln(\mu/\Lambda)$, e_u, e_d are quark charges. The differential operator $1 - M^2\frac{\partial}{\partial M^2}$ is applied to exclude the contributions of $N \to N^*$ transitions. Using the approximate expression for $\tilde{\lambda}_N^2$ from the mass sum rules one can get simple approximated formulae

$$\mu_p = \frac{8}{3}\left(1 + \frac{1}{6}\frac{a}{m^3}\right) = 2.96\ (2.79) \quad (25)$$

$$\mu_n = -\frac{4}{3}\left(1 + \frac{2}{3}\frac{a}{m^3}\right) = -1.93(-1.91)$$

(in parenthesis - experimental numbers). The values of μ_p and μ_n obtained from a more exact treatment of eqs.(23),(24) do not differ from (25) within the expected accuracy of the calculation (±10-15%).

Hyperon magnetic moments were calculated analogously [38]. (Here the strange quark mass m_s and the value f (15) come into play)

The results of the calculation of baryon magnetic moments are given in Table 3 in comparison with the experimental data and with the nonrelativistic quark model calculations.

Table 3.

Baryon octet magnetic moments

	p	n	Σ^+	Σ^0	Σ^-
sum rules	3.0	-2.0	2.4	.70	-1.0
quark model	2.79[b)]	-1.91[b)]	2.67	.78	-1.09
exp	2.79[39]	-1.91[39]	2.38[40] ±.02	-	-1.11[41] ±.03

continued on the next page

Table 3.(continuation)

	Ξ^0	Ξ^-	Λ	$\Lambda \to \Sigma^0$
sum rules	-1.40	-.90	-.70[a)]	1.55[a)]
quark model	-1.44	-.49	-.61[b)]	1.63
exp	-1.25[39] ±.015	-.69[42] ±.04	-.61[39]	$1.82^{+.25}_{-.18}$ [43]

a) Approximate formulae (see [36])

b) Input data

The value of quark condensate magnetic susceptibility χ can be determined from the same sum rules for nucleon magnetic moments as well as by considering a special sum rule [44] for quantity $\chi(q^2)$ defined by

$$\int d^4x\, e^{iqx}\langle 0|T\{\bar{u}(x)\gamma_\lambda u(x), \bar{u}(0)\sigma_{\mu\nu}u(0)\}|0\rangle = (\delta_{\mu\lambda}q_\nu - \delta_{\nu\lambda}q_\mu)\langle 0|\bar{u}u|0\rangle\chi(q^2)$$

It is easy to see that $\chi(0) = \chi$ if the diagrams fig.9b are neglected. For $\chi(q^2)$ on one hand, an operator expansion at large $Q^2 = -q^2$ is written, and on the other hand, the subtractionless dispersion relation which is saturated by ρ and ρ' meson contributions. The parameters of ρ and ρ' are determined from the sum rule obtained. Thus, $\chi(q^2)$ at all $q^2 < 0$ is found. Particularly,

$$\chi = \chi(0) = -5.7 \pm 0.6\ \mathrm{GeV}^{-2} \quad (26)$$

It should be emphasized that VDM leads to twice as least value of χ [37]. Function $\chi(q^2)$ plays an important role in calculating nucleon magnetic formfactors (figs.7,8).

The method similar to that used in the baryon magnetic moments calculations can be applied to calculating the weak interaction constant g_A [45]. To this end one should introduce as an external field the constant isovector axial field A_μ interacting with quarks versus an additional term in lagrangian

$$\Delta L = (\bar{u}\gamma_\mu\gamma_5 u - \bar{d}\gamma_\mu\gamma_5 d)A_\mu \quad (27)$$

An essential simplification appears in this problem since some field induced v.e.v.'s can be determined proceeding from existence of a massless pion. In particular,

$$\langle 0|\bar{u}\gamma_\mu\gamma_5 u|0\rangle_A = f_\pi^2 A_\mu \qquad (28)$$

An important moment in calculating g_A is that the difference of g_A from 1 results only from nonperturbative effects due to quark condensate or due to field induced v.e.v's. As a result, it was obtained [45, 46] that

$$g_A = 1.40 \pm 0.20 \quad (g_A^{exp} = 1.25) \qquad (29)$$

The nucleon interaction constant g_A^s with the axial isoscalar field has been also found [46]:

$$g_A^s = 0.5 \pm 0.2 \qquad (30)$$

The obtained value of g_A^s is of interest in view of the checking the SU(6) symmetric quark model where $(g_A^s)_{SU(6)} = 1$ (recall that $(g_A)_{SU(6)} = 5/3$). Besides, g_A^s enters the Bjorken sum rule for the deep-inelastic scattering of longitudinal polarized electrons on longitudinal polarized nucleons

$$\int_0^1 g_1^{p,n}(x)\,dx = \frac{1}{12}\left(\pm g_A + \frac{5}{3} g_A^s\right) \qquad (31)$$

where the upper and lower signs refer to proton and neutron, respectively. In the exact SU(3) symmetry g_A^s may be expressed via F and D coupling constant in hyperon leptonic decays [47]

$$g_A^s = 3F - D = 0.68 \pm 0.05 \qquad (32)$$

The knowledge of g_A permits one by virtue of the Goldberger-Treiman relation to find the pion-nucleon interaction constant

$$\frac{g^2_{\pi NN}}{4\pi} = \frac{1}{2\pi}\frac{m^2}{f_\pi^2} g_A^2 = 16 \pm 3 \qquad (33)$$

in comparison with $(g^2/4\pi)_{exp} = 14.5$.

The discussed above method of consideration of hadron properties in external fields is quite a general. It is applicable to calculating the emission and absorption amplitudes of the fields with the wavelengths much larger than hadronic dimensions. In particular, one may determine the interaction amplitudes of soft pions with hadrons. In this case an additional term $\Delta L = j^a_{\mu 5} A^a_\mu$ is introduced into lagrangian ($j^a_{\mu 5}$ is the axial isovector current of quarks, A^a_μ is the external axial field considered in the limit when its momentum q tends to zero). The transversality condition of the axial field emission amplitudes $q_\mu T_\mu = 0$ (or $q_\mu T_\mu =$ = terms with current commutators) allows one to express the soft pion emission amplitudes through the interaction amplitudes with the field A_μ in the limit $q \to 0$.

In ref. [49] the sum rules for the vertex function of transition A → B + π (soft) were studied and the interaction constants

$$\frac{g^2_{\pi NN}}{4\pi} = 12.3 \qquad \frac{g^2_{\Sigma\Sigma\pi}}{4\pi} = 10 \qquad \frac{g^2_{\Delta N\pi}}{4\pi} = 18\ \mathrm{GeV}^{-2} \qquad (34)$$

$$g_{\omega\rho\pi} = 12.5\ \mathrm{GeV}^{-1} \qquad (35)$$

(cf. (21))

were found.

Unfortunately, the accuracy of relations (34,35) is unclear since in [49] only the term of the lowest dimension in the operator expansion was taken into account and contributions of nondiagonal transitions $A \to B^*\pi$, $A^* \to B\pi$ were not subtracted.

An interesting application of this method was done in the recent paper [50] where transitions $D^* \to D\gamma$, $D^* \to D\pi$ were considered and the values of partial widths were found. It is important to note that no new unknown condensates enter this calculation. The results are (in keV)

	Br(%)
$\Gamma(D^{*+} \to D^+\gamma) = 0.7 \pm 0.2$	5 (8±7)
$\Gamma(D^{*+} \to D^0\pi^+) = 8.5 \pm 2.5$	66(64±11)
$\Gamma(D^{*+} \to D^+\pi^0) = 3.5 \pm 1$	29(28±9)
$\Gamma(D^{*0} \to D^0\gamma) = 14 \pm 7$	70(45±15)
$\Gamma(D^{*0} \to D^0\pi^0) = 6.0 \pm 2$	30(55±15)

(In the second column the branching ratios are given, in parentheses the experimental data [39]).

It should be, however, borne in mind that when applying this method to nondiagonal transitions $A \to B\gamma$, $A \to B\pi$ it works rather well only when A and B masses are close, $m_A - m_B \ll m_{char} \sim m_\rho$ and A and B are the lowest states in their channels. Otherwise it is difficult to separate the contribution of the transition we are interested in from a large background of transitions into excited states.

4. Structure Functions.

The problem of modelless calculation of structure functions is one of the most important in QCD. Recently it was achieved some progress in this domain: valence quark contribution into structure functions at small x has been calculated [51]. In the calculation it was assumed that at small x quarks distributions had the Regge behaviour

$$u_v(x) - d_v(x) = \beta^{(\rho)}(0)\, x^{-\alpha_\rho(0)}$$
$$u_v(x) + d_v(x) = \beta^{(\omega)}(0)\, x^{-\alpha_\omega(0)} \quad (36)$$

Intercepts $\alpha_\rho(0) = \alpha_\omega(0) \approx 1/2$ were taken from experiment and coefficients $\beta^{(\rho)}(0)$ and $\beta^{(\omega)}(0)$ were determined. The idea of the calculation was based on QCD sum rules consideration of the upper vertex in the graph of fig.10 near $t=(q'-q)^2 = m_\rho^2$ and afterwards in extrapolation of the amplitude from $t=m_\rho^2$ into t=0.

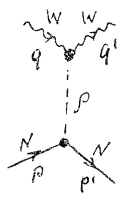

Fig.10.

The resultant expression for small x-distribution of valence quarks in proton is (at intermediate $Q^2 \sim$ 5-20 GeV^2)

$$u_v(x) = 1.65x^{-1/2} \qquad d_v(x) = 0.83x^{-1/2} \quad (37)$$

and in pion

$$u_v^{\pi^+}(x) = 0.68x^{-1/2} \quad (38)$$

(The accuracy of the coefficient at $x^{-1/2}$ is about 30%). The theoretical expectations (37),(38) are in agreement with experiment.

5. Numerical Values of Vacuum Condensates.

As was shown above the operator expansion method - QCD sum rule method - is able to explain and to predict a lot of observable phenomena using a very little number of phenomenological constants - vacuum condensates. Numerical values of these constants are extremely important - the above considered physics is based on it. Therefore it is worth dwelling on the consideration of this problem. When considering baryons the most essential is the value of quark condensate $\langle 0|\bar{q}q|0\rangle$, q=u,d. Its low-energy limit is determined by the well known formula

$$\langle 0|\bar{q}q|0\rangle = -\frac{1}{2}\frac{m_\pi^2 f_\pi^2}{m_u+m_d}, \quad f_\pi = 133\,MeV \quad (39)$$

The thorough analysis done by Gasser and Leutwyler [52] taking into account the first terms of expansion in m_s showed that

$$2\frac{m_s}{m_u+m_d} = 25.7 \pm 2.6 \quad (40)$$

I take for m_s the value found from the baryon mass best fitting $m_s = 150 \pm 30$ MeV [6,38] (at normalization point μ =0.5 GeV). This value agrees with the value found recently from the best fit of the vector meson sum rules [53] m_s(0.5 GeV) = 130±12 MeV (although the error quoted in ref. [53] seems to me underestimated). Then I find (at μ = 0.5 GeV)

$$\langle 0|\bar{q}q|0\rangle = -(240\ MeV)^3 \pm 25\% \quad (41)$$

If Λ_{QCD} is taken to be 100 MeV, then the renorminvariant quantity

$$\alpha_s \langle 0|\bar{q}q|0\rangle^2 = (0.8^{+0.6}_{-0.4}) \cdot 10^{-4} GeV^6 \quad (42)$$

The numerical value of the 5-dimensional quark-gluon condensate

$$-g\langle 0|\bar{q}\sigma_{\mu\nu}\frac{\lambda^n}{2}G^n_{\mu\nu}q|0\rangle \equiv m_0^2\langle 0|\bar{q}q|0\rangle \quad (43)$$

was determined from selfconsistency conditions of baryon sum rules [6]. The result is m_0^2 = (0.8±0.3) GeV^2.

It seems to me that the most plausible value of $f = \langle 0|\bar{s}s|0\rangle / \langle 0|\bar{u}u|0\rangle - 1$ is given above (see eq.(15)) though in the literature there are also other estimate [54-56].

The value of gluon condensate $\langle 0|(\alpha_s/\pi)G^n_{\mu\nu}G^n_{\mu\nu}|0\rangle = 0.012\ GeV^4$ found in ref. [1] was taken in the ITEP calculations. It is possible that its real value is by 30-40% larger. A considerably larger increase of gluon condensate proposed in ref. [57] is rejected by the analysis of entire data [58] and is not needed for agreement of charmonium sum rules [59].

Now about the factorization hypothesis. Though the proof is absent (besides the one in the limit $N_c \to \infty$) it is a common belief that for 4-quarks v.e.v.'s the factorization hypothesis works well (with an accuracy of about 10%). For 4-gluon v.e.v's

$\langle 0|G^4|0\rangle$ there are doubts if the factorization works, especially for those which are nonleading at $N_c \to \infty$ (proportional to N_c^3 instead of N_c^4)[58]. The estimate done in ref.[60] show that the contribution of the most important in charmonium sum rule calculation 4-gluon operator[57] is probably overestimated by the use of factorization hypothesis. For v.e.v's of higher dimensions the situation is even worse: the results depend on which way the factorization hypothesis is applied (for example, before x-expansion or after it).

6. Heavy Quarkonium

The great success of QCD in describing heavy quarkonium levels is well known (see, for example, reviews[61,62]). In such a description an important role is played by nonperturbative interaction of heavy quarks with long-wave vacuum fluctuations of the chromoelectric gluonic field which results in appearance of gluonic condensate $\langle 0|(\alpha_s/\pi)\, G^n_{\mu\nu} G^n_{\mu\nu}|0\rangle$ in the formulae for energy levels.

I will not discuss this point here but turn to the problem where a certain understanding arose last time, namely to the problem of nonperturbative effects in hadronic annihilation of heavy quarkonia. According to the celebrated Appelquist-Politzer recipe[63] it was usually accepted that the total widths of direct decays $(\bar{Q}Q) \to$ hadrons can be determined calculating the $(\bar{Q}Q) \to 2$ gluon widths for C-even quarkonium states or $(\bar{Q}Q) \to 3$ gluons for C-odd ones. In doing so, the ratio of direct (not via the virtual photon) decay $J/\psi, \psi', \ldots, \Upsilon, \Upsilon', \ldots$ into hadrons to the width of their decay into $\mu^+\mu^-$ is given by the formula

$$R_h = \frac{\Gamma(^3S_1 \to \text{direct hadrons})}{\Gamma(^3S_1 \to \mu^+\mu^-)} = \frac{\Gamma(^3S_1 \to 3g)}{\Gamma(^3S_1 \to \mu^+\mu^-)} = \frac{10}{81}\,\frac{\pi^2-9}{\pi}\,\frac{\alpha_s^3(m_Q)}{\alpha^2 e_Q}\left(1+\frac{\alpha_s}{\pi}\right) \qquad (44)$$

where correction $\sim\alpha_s$ is calculated in the renormalization scheme $\overline{\text{MS}}$[49]. By comparing $\Gamma(J/\psi \to 3g)/\Gamma(J/\psi \to \mu^+\mu^-)$ with experiment it was originally found $\alpha_s(m_c) = 0.2$ and was deduced that $\Lambda_{QCD} = 100$ MeV[1]. But later, basing on a more wide and exact set of experimental data it became clear that the Appelquist-Politzer recipe does not result to rather good accuracy prediction for hadronic width of J/ψ. The arguments are the following. Let us suppose that the Appelquist-Politzer recipe works well for upsilonium. From experimental value $R_h(\Upsilon)$[65,109] we can find $\alpha_s(m_b) = 0.165 \pm 0.005$ (see[65]), determine by renormalization group equations $\alpha_s(m_c)$ and calculate according to (44) $R_h(J/\psi)$. Then the obtained value will be more than by 1.5 times larger than the experimental one. The appeared discrepances can be only due to nonperturbative effects in $\Gamma(^3S_1(\bar{Q}Q) \to \text{hadrons})$. The nonperturbative effects in very heavy quarkonium when the levels are almost Coulomb have been considered by Voloshin[66]. The initial point of his approach is the studying of imaginary part of the $\bar{Q}Q$ system Green function under conditions when $k_n R \gg 1$ where k_n is the heavy quark momentum in quarkonium, $k_n = (2/3)\,\alpha_s(k_n)\, m_Q/n$, n is the principal quantum number, R is the confinement radius. In this case the heavy quark interaction with long-wave fluctutations of gluonic field can be taken into account making use of multipole expansion and taking only chromoelectric (EI) and chromomagnetic (MI) terms. When averaging gluonic field over fluctuations the linear in field terms drop out and the effect appears in the second order in the field, i.e. turns out to be proportional to $\langle 0|(\alpha_s/\pi)\, G^n_{\mu\nu} G^n_{\mu\nu}|0\rangle$. Nonperturbative corrections to ratio R_h (44) are described by the graphs of fig.11a,b where wide bright lines denote the Green function of $\bar{Q}Q$ pair in colourless state, thick lines - the Green functions of $\bar{Q}Q$ pair in the octet colour state, dashed lines - are hard gluon propagators. (The contribution of graphs where both the interactions with vacuum fluctuations occur before or after annihilation drops out in the ratio R_h).

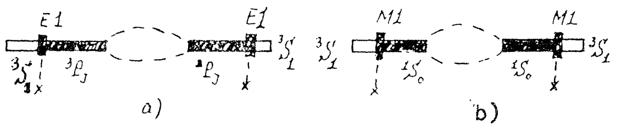

Fig.11.

Though MI interaction contains a small parameter v/c, the contributions of MI and EI interactions appear to be of the same order since the latter converts the initial S-wave into P-wave whose annihilation amplitude is suppressed by the factor v/c. The nonperturbative correction to the ratio

R_h (44) for the n-th level of quarkonium calculated by Voloshin [66] turned out to be

$$\frac{\delta R_h}{R_h}(n^3S_1) = -\frac{\pi^2 a_n \langle 0|(\alpha_s/\pi)\, G^a_{\mu\nu} G^a_{\mu\nu}|0\rangle}{2^9 \cdot 9(\pi^2-9)\alpha_s(m_Q)\, k_n^4 n^4} \qquad (45)$$

where a_n are numerical coefficients, $a_1 = 4$, $a_2 = 9$ (a_n, $n \geq 3$ are very small). Substituting into (45) the numbers for Υ ($k_1 = 0.96$ GeV) and for J/ψ ($k_1 = 0.5$ GeV), we find $\delta R_h/R_h(\Upsilon) = 0.4\%$ and $\delta R_h/R_h(J/\psi) = 5\%$. It should be, however, borne in mind that in a_1 a strong (by the order of magnitude) compensation of EI and MI transition contributions arises which is caused by the form of the quarkonium Coulomb wave function. Since Υ , and, moreover, J/ψ states have nothing in common with Coulomb ones, extrapolation of relations (45) to these states should be done with a care and it is not excluded that in J/ψ the true nonperturbative effect is several times larger than that given by formula (45). Notice that the sign of the effect in (45) coincides with that required by experiment for J/ψ

Since the nonperturbative correction (45) for Υ is very small there is a confidence that with an accuracy not worse than a few per cent the ratio $R_h(\Upsilon)$ is given by formula (44) and consequently determination of $\alpha_s(m_b)$ from here is sufficiently reliable. Exact measurements of the ratio $R_h(\Upsilon')/R_h(\Upsilon)$ (it is expected that it is equal to 1 up to a few per cent) could confirm or deny this statement. (The present experimental value of $R_h(\Upsilon')/R_h(\Upsilon)$ is 1.11 ± 0.25 [109]). Other good method for finding $\alpha_s(m_b)$ is the measurement of the ratio [67] $R_\gamma(\Upsilon) = \Gamma(\Upsilon \to \gamma + hadr)/\Gamma(\Upsilon \to hadr)$

$$R_\gamma(\Upsilon) = \frac{36}{5} e_b^2 \frac{\alpha}{\alpha_s(0.27 M_\Upsilon)} \left[1 + (0 \pm 0.6)\frac{\alpha_s}{\pi}\right] \qquad (46)$$

where nonperturbative correction is by a factor of three smaller than formula (45).

Nonperturbative corrections to hadronic widths of quarkonium P-levels can be calculated analogously. They appear to be much more than for S-levels; even for b-quarks they comprise few dozens per cent. Thereby, perturbation theory calculations of $\bar{b}b$-quarkonium P-level hadronic widths have a bad accuracy and for charmonium they are not confident at all.

II. Perturbative QCD.

1. Multiplicity and Inclusive Spectra in Hard Processes.

I will discuss only those problems of perturbative QCD where a noticeable progress has been achieved in the last few years. First, it refers to the problem of multiplicity and inclusive spectra in jets in hard processes, in particular, to the role of coherent effects in these phenomena [68-77]

I start with considering the space-time picture of development of a quark jet in e^+e^- annihilation at high energy E of colliding beams [77] . It can be shown from uncertainty relation $\Delta E \Delta t \sim 1$ that the characteristic time t which is necessary for the fast quark emission of gluons with momentum components $k_\parallel$, $k_\perp$ is of order (in lab. system, $E \gg k_\parallel \gg k_\perp$)

$$t \sim \frac{k_\parallel}{k_\perp^2} \qquad (47)$$

i.e. is very large at $k_\parallel \gg k_\perp$. In the quark rest system the characteristic momenta of a related to quark gluonic cloud are $\tilde{k}_\perp \sim \tilde{k}_\parallel \sim R^{-1}$ where R is the confinement radius. In laboratory system we have

$$k_\perp = \tilde{k}_\perp \sim R^{-1}, \quad k_\parallel = \frac{E}{m}\tilde{k}_\parallel \sim \frac{E}{m} R^{-1} \qquad (48)$$

and consequently the time needed for a fast quark to create a surrounded gluonic cloud is

$$t \sim \frac{E}{m} R \qquad (49)$$

(For light quarks in (49) m should be implied as $\langle k_\perp^2 \rangle^{1/2}$ or R^{-1}). Therefore, at $E/m \gg 1$ the quark strong interaction with gluonic field - formation of a string between a $\bar{q}q$ pair could occur only at very large distances between components of the pair, which, however, contradicts causality. Hence it follows that in e^+e^- annihilation (as well as in other hard processes) long strings between quarks (or gluons) cannot be produced. Thus, at initial stage of e^+e^- annihilation (or of any other hard processes) multiplication can occur only through bremsstrahlung gluon emission which featured by a small coupling constant $\alpha_s(k_\perp^2)$.

It is seen from relation (47), if one re-writes it as

$$t \sim \frac{1}{k_\perp \theta} \qquad (50)$$

where θ is the bremsstrahlung gluon emission angle, that small t $1/E<t\ll R$ correspond to $E>k_\perp\theta\gg R^{-1}$. The account of small angle θ region and of a wide interval of transverse momenta $E\gg k_\perp\gg R^{-1}$ leads to appearance of the double logarithmic factors $\ln^2 E/\Lambda$ compensating the smallness of α_s. in the gluon emission probability. It can be shown that in considering parton (bremsstrahlung gluons and quark pairs) emission in the double logarithmic approximation subsequent partons emit subsequently in time, $t_{n+1} > t_n$. This makes it possible to probabilistically describe the multiplication process in jets. In such a manner the heavy quark pair multiplicity in jets was originally calculated in refs.[78,118] The further studying has demonstrated, however, that the method of refs[78,118] is invalid since it does not take into account coherence effects in gluon emission. The physics here is very simple and can be explained on an example of quantum electrodynamics where the existence of such a phenomena was originally pointed out by Chudakov[79]. Consider the electron-positron high energy pair production and photon emission by this pair. It is clear that the photon with transverse momentum relative to pair motion $k_\perp$ can emit only if $k_\perp^{-1}<\rho_\perp$ where $\rho_\perp$ is the transverse distance between components of the pair. (In the opposite case $k_\perp^{-1}>\rho_\perp$ there is a destructive interference of e^+ and e^- emission). The value of $\rho_\perp$ can be estimated as $\rho_\perp\sim t\,\theta_0$ where θ_0 is the angle between e^+ and e^- momenta and t is the formation time of the emission (50). Hence it follows that

$$\theta_0 > \theta \qquad (51)$$

This means that the angles of subsequent emissions are ordered[68, 69]. Thus, to obtain correct formulae in the double-logarithmic approximation one should take into account the angular ordering of subsequent emissions while in the rest one may use the former classical probabilistic treatment. Qualitatively, changes in comparison with the former results of ref.[78] are evident: since the probabilistic description preserves, cross sections become smaller and owing to conditions $k_\perp = Ex\,\theta > Q_0$ such a decrease is to refer to the small x region - a dip in the small rapidity region must appear in the rapidity distribution.

The discussed above coherence effects are of a very general character and manifest themselves in any hard processes. Another approach for consideration of these effects was developed in ref.[74] where a generalized coherent state operator was constructed and the Bloch-Nordsiek type cancellation in inclusive processes ($e^+e^-\to$ hadrons, deep inelastic scattering and Drell-Yan process) was proved.

The results obtained in the double logarithmic approximation[71-73] give the correct qualitative picture but they are inapplicable for quantitative predictions since single-logarithmic corrections are very large (their smallness is only $\sqrt{\alpha_s(E)}$) and in a number of cases lead to pre-exponential factors chaning even the asymptotic behaviour.

Mueller[75] and Dokshitzer and Troyan[76] have forwarded an idea how to take into account single-logarithmic terms. According to hypothesis[76] confirmed by perturbation theory calculations up to the 4-th order the A $\to$ BC decay probability (A,B,C are gluons G and quarks q_f)

$$dw=\frac{\alpha_s(k_\perp)}{2\pi}\Phi_A^{BC}(z)V(\vec n)dz\,\frac{d\Omega}{8\pi} \qquad (52)$$

where $\Phi_A^{BC}(z)$ are common kernels of the evolution equations

$$V^s_{f(g)}(\vec n)=\frac{a_{sg}+a_{fg}-a_{sf}}{a_{sf}\,a_{sg}},\quad a_{ik}=1-\cos\theta_{ik} \qquad (53)$$

and indeces refer to the emitted gluon - the "son" (s) and its "father" (f) and "grandfather" (g). The angular kernel (53) has the property

$$\int\frac{d\varphi_{s(f)}}{2\pi}V^s_{f(g)}(\vec n)=\frac{2}{a_{fg}}\,\theta(a_{fg}-a_{sf}) \qquad (54)$$

i.e. typical of the double-logarithmic approximation condition $\theta_{sf}\ll\theta_{fg}$ is replaced by strict inequality $\theta_{sf}<\theta_{fg}$. Making use of (52) and preserving the probabilistic treating of parton multiplication in jets the authors of[76] managed to determine inclusive spectra of partons in jets. Similar formulae but with another technique have been obtained by Mueller[75]. The soft gluon distribution in a jet produced by a hard gluon or a quark was found. From this distribution the multiplcity of gluons with energies $E_g > Q_0$ in a gluonic jet was calculated:

$$n_g^g = z\left(\frac{z}{z_0}\right)^{-B}\left[I_{B+1}(z)K_B(z_0)+K_{B+1}(z)I_B(z_0)\right] \quad (55)$$

where $z=\sqrt{(16N_c/b)Y}$, $z_0=\sqrt{(16N_c/b)\lambda}$
$Y=\ln(E/\Lambda_{QCD})$, $\lambda=\ln(Q_0/\Lambda_{QCD})$,
$b=(11/3)N_c-(2/3)n_f$, $B=[(11/3)N_c+(2/3)(n_f/N_c^2)]/b$, I,K are the Bessel functions from imaginary argument, N_c is the colour number, n_f is the flavour number. The gluon multiplicity in the quark jet $n_q^g \approx (C_F/N_c)\,n_g^g = (4/9)\,n_g^g$
In the asymptotics at $E/Q_0 \to \infty$

$$n_g^g \sim \left(\ln\frac{E}{\Lambda}\right)^{-\frac{B}{2}+\frac{1}{4}} \exp\left[\sqrt{\frac{16N_c}{b}\ln\frac{E}{\Lambda}}\right] \quad (56)$$

To arrive from here at physical conclusions one has to admit some hypothesis about the hadronization mechanism, i.e. how gluons and quarks transform into physical hadrons. The soft screening hypothesis when gluon with virtuality $Q_0 \sim 1$ GeV goes over into finite number of hadrons seems to be natural [77]. The arguments in favour of such a hypothesis is quark-hadron duality as well as experimentally observed adjacent correlation in rapidity spectra. The gluon inclusive spectrum calculated in the double+single- logarithmic approximation tends to finite limit at $Q_0 \to 0$, i.e. the necessary condition for this hypothesis to work is fulfilled. An attempt was made in ref. [77] to describe particle spectra observed in e^+e^- annihilation assuming the finite gluon to go over into K_h hadrons of the type h, K_h = Const. Fig.12 shows the pion spectrum obtained in [77] in comparison with experimental data at 2E=34 GeV [80].

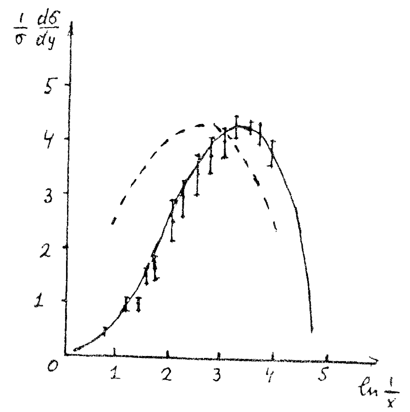

Fig.12

(K_h = 1.09, it was put in the calculation that $Q_0 \simeq \Lambda \approx 150$ MeV). The experimental data reveal a trend to formation of a dip at small x though final conclusions would be premature.

Notice that the spectrum in the double-logarithmic approximation without taking into account one-logarithmic terms appears to be quite different (dotted line in fig. 12).

New experimental data presented at this conference give further support to the existence of coherence effects in multiparticle production processes. The rapidity distributions of π and K in e^+e^--annihilation measured by the PEP-4-LBL collaboration [110] at 2 E = 29 GeV shows a clear dip at small y. An interesting confirmation of coherence effects was found by the JADE collaboration [111] in the study of angular dependence of hadronic multiplicity in three (quark-antiquark-gluon) jet events in e^+e^--annihilation. It was found that the multiplicity in the region between q and $\bar{q}$ jets is smaller than between q and g or $\bar{q}$ and g jets. This fact is a simple consequence of destructive interference of bremsstrahlung gluon radiation emitted by quark and antiquark and can be easily understood by drawing the colour line (so called t'Hooft) diagrams.

In the papers by Marchesini and Webber [81] and Webber [82] another approach to the consideration of the problem was proposed. This approach is grounded on the above physical principles, but exploit the numerical Monte-Carlo method of calculation. Since strict inequalities on the angles of subsequent gluon (or quark) emissions $\theta_{sf} < \theta_{fq}$ were admitted in the Monte-Carlo calculations this method automatically corresponds to the double + single logarithmic approximation. The results of both approaches are in agreement with one another, but the Monte-Carlo method has the advantage that its results can be compared with experiment in a more direct way. Such a comparison was performed by the JADE group [111] and good agreement between theory and experiment was found.

In the paper [112] presented to this conference the angular ordering (in double logarithmic approximation) was proved using only the properties of unitarity and gauge invariance without laborious analysis of high order Feynman graphs. Angular ordering was obtained as a condition for the separation of the loop with the softest particle from the amplitude. After that the fragmentation function and parton multiplicity in the jet was found in a rather simple way.

2. Heavy Quark Production

Production of hadrons containing a heavy quark Q (of mesons $H=\bar{q}Q$, baryons $B=Qqq$) is of special interest since this is for the present the only case when produced hadron distribution can be calculated with perturbative QCD. The production of heavy quarks is also interesting for another reason: the experimental study of jet asymmetry in the process $e^+e^- \to \bar{Q}Q$ at high energies opens a new way in investigation of weak interactions. In the general case we can write for inclusive spectrum of heavy mesons H as

$$\mathcal{D}_Q^H(x) = \int_x^1 \frac{dx_Q}{x_Q} \mathcal{D}_Q(x_Q)\, w_Q^H(x/x_Q) \qquad (57)$$

where $\mathcal{D}_Q(x_Q)$ is the probability for quark Q to have a fraction of momentum x_Q as a result of hard gluon emission, $w_Q^H(x/x_Q)$ is the probability of quark fragmentation into meson H. For very heavy quark $w_Q^H(x) \approx \delta[1-x-O(1/MR)]$ where M is the quark mass, R is the confinement radius [83 - 85]. If energy E is not very high $[\alpha_s(M^2)/\pi]\, \ln(E^2/M^2) << 1$ then $\mathcal{D}_Q(x_Q)$ is determined by single gluon emission [86]. At very high E $D_Q(x)$ was calculated in the leading logarithmic approximation and can be represented by formula [87]

$$\mathcal{D}_Q(x) = A\, \frac{1+x^2}{2} (1-x)^{-1+(16/3)\xi}$$

$$A^{-1} = \frac{3}{16\xi}\left(1+\frac{16}{3}\xi\right)^{-1} + \frac{1}{4}\left(1+\frac{8}{3}\xi\right)^{-1}, \quad \xi = \frac{1}{6}\ln\frac{\alpha_s(M^2)}{\alpha_s(E^2)} \qquad (58)$$

The average fraction of the energy shared by quark Q is equal (in the leading and next-to-leading logarithmic approximation): to :

$$\langle x_Q \rangle = \left[\exp\left(-\frac{32}{9}\xi\right)\right]\left(1+\frac{44}{27}\frac{\alpha_s}{\pi}\right) \qquad (59)$$

At not very small μ/M function w_Q^H differs from δ -function and is model-dependent. It seems that when H is a boson, the most appropriate form is the suggested in [88] expression based on the quantum-mechanical relation $w_Q^H \sim 1/(\Delta E)^2$ $\Delta E = E_Q - E_H - E_q$ Then

$$w_Q^H(x) = \frac{N}{x}\left[-1+\frac{1}{x}+\frac{\varepsilon_Q}{1-x}\right]^{-2} \qquad (60)$$

where N is the normalization factor determined from condition

$$\int_0^1 w_Q^H(x)\, dx = 1$$

and ε_Q is unknown parameter, $\varepsilon_Q \sim (\mu/M)^2$.

The above presented formulae can be used in analysing the charmed B-meson production in e^+e^- annihilation. The c-quark mass is not large enough for quantitative comparison of these relations with experiment: nonperturbative effects here are still uncontrollable.

As follows from (57) the mean momentum carried by heavy meson is equal

$$\langle x \rangle = \langle x_Q \rangle \langle z \rangle$$

where

$$\langle z \rangle = \int_0^1 z\, w_Q^H(z)\, dz$$

Numericaly, for B meson (at $\sqrt{\varepsilon} = 0.06$) we have $\langle x_\varepsilon \rangle_{theor} = 0.80$ in good agreement with experimental data [110,119] $\langle x_\varepsilon \rangle_{exp} = 0.80 \pm 0.10$.

The heavy meson (B-meson) production process in the region $1-x >> \mu/M$ is of great interest also from another point of view - as a trigger on a "pure" gluonic jet [86]. (At accessible energies B-meson production with $1-x >> \mu/M$ is mostly accompanied with one gluon emission).

3. Exclusive Processes.

In the last few years a great progress has been achieved in perturbative QCD description of exclusive processes with large momentum transfers, such as formfactors, large angle cross sections etc. (see, e.g. review [31]). The results obtained endure various theoretical tests and it seems, there are no grounds for doubts about their reliability in the asymptotic region $Q^2 \to \infty$. However, as it was shown in the first part of my talk, in the region of not very large Q^2 practically accessible on modern accelerators these results contradict experiment and QCD sum rule calculations as well as some model calculations. Such contradictions emerge in formfactors of pion, of ρ-meson, of $\rho \to \pi$ transition (at $Q^2 \sim 5$ GeV^2), in the nucleon formfactor (at $Q^2 \sim 25$ GeV^2). The α_s corrections do not improve the situation [89]. Therefore, there is an impression that the applicability region of exclusive process perturbative theory is achieved at very large Q^2. As a consequence a question arises what is the origin of this effect and is it possible to improve the theory in such a way that its applicability region will begin far earlier. An attempt to answer this question was made in the works by Chernyak and Zhitnitsky [90]. I explain the idea on an example of the pion formfactor. At large Q^2 the the pion formfactor can be written as [21-25]

$$F_\pi(Q^2) = \frac{32\pi}{9}\frac{\alpha_s(Q^2)}{Q^2} f_\pi^2 \int_{-1}^{1} d\xi \int_{-1}^{1} d\xi' \frac{1}{(1-\xi)(1-\xi')} \times \varphi(\xi, Q^2)\varphi(\xi', Q^2) \tag{61}$$

where $\varphi(\xi, Q^2)$ has the meaning of the wave function of $u\bar{d}$ quarks in the pion, ξ is the difference of u- and d-quark momenta (in units of the total pion momentum in the Breit system). $\varphi(\xi, Q^2)$ is formally defined by equation

$$\langle 0|\bar{d}(z)\gamma_\mu\gamma_5 \exp\left[ig\frac{\lambda^n}{2}\int_{-z}^{z} dx_\mu A_\mu^n(x)\right]u(-z)|\pi^+(p)\rangle = ip_\mu f_\pi \int_{-1}^{1} d\xi\, \varphi(\xi, z^{-2}) e^{i\xi(zp)} + \text{terms} \sim z_\mu \tag{62}$$

In the leading logarithmic approximation it can be put in (62) $z^2 \sim 1/Q^2$. In accord with (62) $\varphi(\xi, Q^2)$ is normalized by condition

$$\int_{-1}^{1} \varphi(\xi, Q^2) d\xi \Big|_{Q^2\to\infty} = 1 \tag{63}$$

The usual method of treating relations (61),(62) is to expand $\varphi(\xi, Q^2)$ in orthogonal polynomials chosen in such a way that expansion coefficients $\varphi_n(Q^2)$ transform multiplicatively at renormalization group transformations. In the case of the pion formfactor they are Gegenbauer polynomials $G_{2n}^{3/2}$ and

$$\varphi_n(\xi, Q^2) = \frac{3}{4}\sum_n (1-\xi^2)\varphi_n(Q^2)\, G_{2n}^{3/2}(\xi) \tag{64}$$

At large Q^2 $\varphi_n(Q^2) = [\alpha_s(Q_0^2)/\alpha_s(Q^2)]^{-\gamma_n/6}\varphi_n(Q_0^2)$

Since γ_n increase with n, $\gamma_{n+1} > \gamma_n$, then in the asymptotics in series (64) only the first term survives and

$$\varphi_{asymp}(\xi) = \frac{3}{4}(1-\xi^2) \tag{65}$$

Substituting (65) into (61) gives the asymptotic formula for formfactor $F_\pi(Q^2)$. It is clear, however, that if there is no numerical smallness of higher coefficients, the series in (64) converges very slowly and at accessible Q^2 the replacing $\varphi(\xi)$ by $\varphi_{asymp}(\xi)$ is not correct.

Chernyak and Zhitnitsky [90] tried to construct $\varphi(\xi, Q^2)$ in the region of not very large Q^2 using the information on momenta

$$\langle \xi^n \rangle = \int_{-1}^{1} d\xi \cdot \xi^n \varphi(\xi) \tag{66}$$

which can be obtained with the help of QCD sum rules. As follows from (62), momenta $\langle \xi^n \rangle$ are determined by the matrix elements

$$\langle 0|\bar{d}\gamma_\mu\gamma_5(iz\overleftrightarrow{\nabla})^n u|\pi(p)\rangle, \quad \overleftrightarrow{\nabla} = \overrightarrow{\nabla} - \overleftarrow{\nabla}$$

which can be found by considering the polarization operator

$$\Pi_{\mu\nu}^{(n)} = i\int d^4x\, e^{iqx}\langle 0|T\{\bar{d}(x)\gamma_\mu\gamma_5(iz\overleftrightarrow{\nabla})^n \times u(x), \bar{u}(0)\gamma_\nu\gamma_5 d(0)\}|0\rangle \tag{67}$$

if one applies the sum rule technique to it. It is clear a priori that the account in (67) of the simplest quark loop only corresponds to asymptotic values of momenta, i.e. to the asymptotic wave function $\varphi_{asymp}(\xi)$ (65). Hence it follows the qualitative conclusion: the sign and scale of deviation of $\langle \xi^n \rangle$ from $\langle \xi^n \rangle_{asymp}$ are determined by the power corrections in the operator expansion of $\Pi_{\mu\nu}^{(n)}(q)$. Namely, if power corrections have the same sign as quark loop, then $\langle \xi^n \rangle > \langle \xi^n \rangle_{asymp}$ and the wave function $\varphi(\xi)$ is wider than $\varphi(\xi)_{as}$ (such a situation takes place in the pion case), if the sign of power corrections is opposite to the sign of the quark loop, then $\langle \xi^n \rangle < \langle \xi_n \rangle_{asymp}$ and $\varphi(\xi)$ is narrower than $\varphi(\xi)_{as}$ (the example is the wave function of transverse-polarized ρ-meson). In papers [90] the 2-d and the 4-th momenta of the pion wave function have been found, $\langle \xi^2 \rangle_{\mu_0} = 0.46$ and $\langle \xi^4 \rangle_{\mu_0} = 0.30$ ($\mu_0 = 0.5$ GeV is the normalization point) and the corresponding to it form of the wave function $\varphi_\pi(\xi) = (15/4)(1-\xi^2)\xi^2$ has been proposed. Analogous consideration was made for ρ-meson. This allowed the authors to calculate a great number of characteristics of exclusive processes (of formfactors, charmonium decay branching rates etc.). It should be, however, borne in mind that high momenta cannot be obtained with this method: the region where the power corrections and higher state contribution are both small shrinks with increasing n and even for n=4 the result is not confident. Though the difference of the 2-d momentum from $\langle \xi^2 \rangle_{asym} = .20$ does not give rise to doubt, the accuracy in $\langle \xi^2 \rangle$ is not completely clear because of uncertainty in excited state contribution. The same remark refers to the found in [90] wave functions of other particles.

In case of proton, of interest is the obtained with the same method in [90] conclusion that the fraction of proton momentum carried by u-quark with spin forwarded along the proton spin is considerbaly larger than that carried by each of the rest quarks.

Though the improvement suggested in [90] makes it possible to hope for advancing the perturbative description of exclusive processes into the smaller Q^2 region, it is unlikely for it to work at Q^2 smaller than 10-20 GeV^2 for mesons and than 25-50 GeV^2 for baryons. The basis of this point of view is that, as was mentioned above, meson formfactors up to $Q^2 \sim 3$ GeV^2 are well described by quite another nonperturbative mechanism. Possibly, in the intermediate region $Q^2 \sim 10$-50 GeV^2 of importance are preasymptotic terms corresponding to higher twists. As was shown in [91], such terms give a large contribution into the pion formfactor. In report [117] presented at this conference, a formal method of construction of hadron operators on the light cone was suggested, which may be useful for calculation of asymptotic hadronic wave functions.

4. Other Problems: $\gamma\gamma$ Scattering, Factorization in the Drell-Yan Process etc.

The theoretical situation in $\gamma\gamma$ - scattering has been discussed in sufficient detail in the talk by Glück at the Brighton Conference. The main conclusion is the following [92, 93]. Despite primary hopes [94] the photon structure function $F_2^\gamma(x, Q^2)$ even at very large Q^2 essentially depends on the given at smaller $Q^2=Q_0^2$ hadronic contribution. Therefore, parameter Λ_{QCD} can be determined only from comparing with experiment the evolution of $F_2^\gamma(x, Q^2)$ as a function of Q^2. Thus, the situation with determining Λ from $\gamma\gamma$- scattering is not a bit better than with determining Λ from the deep inelastic eN-scattering.

The factorization problem in the Drell-Yan process has been faced at perturbative QCD for a long time. It can be formulated as follows. In parton model the production cross section of a lepton pair with a mass M in collision of two hadrons A and B is written as

$$\frac{d\sigma}{dM^2dx} = \frac{4\pi\alpha^2}{9M^2}\sum_a{}' \int_0^1 dx_a \int_0^1 dx_{\bar{a}} \left[\mathcal{D}_{a/A}(x_a, M^2)\mathcal{D}_{\bar{a}/B}(x_{\bar{a}}, M^2) + \mathcal{D}_{\bar{a}/A}(x_{\bar{a}}, M^2)\mathcal{D}_{a/B}(x_a, M^2)\right]\delta(x_a x_{\bar{a}} s - M^2)\delta(x + x_{\bar{a}} - x_a) \quad (68)$$

where $\mathcal{D}_{a/A}(x_a, M^2)$, $\mathcal{D}_{\bar{a}/B}(x_{\bar{a}}, M^2)$ are the densities of partons a and $\bar{a}$ in hadrons A and B, $x=E_q/E$ is the ratio of the pair momentum to particle A momentum in c.m.s. When proving (68) in QCD corrections to this relation appear by virtue of exchange between target and emitting particle by soft and collinear gluons which lead to infrared divergences. In order for (68) to take place it is necessary to prove, at least in the leading terms in $1/M^2$ that these corrections either cancel or reduce to renormalization of structure functions $\mathcal{D}_{a/A}$, $\mathcal{D}_{\bar{a}/B}$. In 1981 Bodwin et al. [95] call in question the validity of relation (68) pointing out the graphs corresponding to interaction in initial state of the emitting active quark with the quark-spectator of the target (and vice versa), which in non-abelian theory might result in violation of eq.(68). A detailed analysis made in the one- [96] and two-loop [97] approximation showed that in fact such divergences cancelled. In the recent paper by Collins, Soper and Sterman [98] it is argued that factorization is true in all orders of perturbation theory. Thereby, I hope this problem comes to a happy end. In the Drell-Yan process at accessible energies there are large nonleading logarithmic corrections to the double logarithmic formulae found in the pioneering work by Dokshitzer, D'yakonov and Troyan [99]. These corrections are especially important in the region of small transverse momentum of the leptonic pair $q_\perp \ll M$. An algorithm to calculate the corrections order by order was proposed in ref [100, 114] and direct calculations of α_s and α_s^2 corrections were performed in refs. [101,102,114-116]. It was shown [116], that α_s corrections are important, but at $M > 10$ GeV the α_s^2 correction term is already small and can be neglected. As a result a reliable prediction was obtained for the production of W and Z bosons at the SPS-collider in the region $3\ GeV < q_\perp << M_{W,Z}$, where nonperturbative corrections as well as higher order perturbative corrections are rather small.

The long standing problem of calculating higher twist corrections in deep-inelastic processes receives its further development in paper [103] where an attempt was made to find leading power-law corrections in the region $x \to 1$. The authors used and extended the technique first employed by Berger and Brodsky [104] in the calculations of power corrections to pion structure function. For eN-scattering they found a very large, proportional to $1/Q^2$ correction to ratio σ_L/σ_T and a substantional one $\sim (1-x)^2/Q^2$ to νW_2. However, the results cardinally depend on the proposed form of the nucleon wave function and therefore its reliability is not clear.

The existence and proposed in ref. [104] form of twist four corrections to the pion structure function was confirmed by the new data presented at this Conference [120,121].

The problem of determination of strong coupling constant α_s (or QCD parameter Λ_{QCD}) from experimental data is intensively discussed for many years. As is well known the value of $\alpha_s(Q^2)$ at fixed scale Q^2 depends on the choice of renormalization scheme. The most popular from these schemes is so called $\overline{MS}$ scheme and usually the values of α_s and Λ are given in this scheme. The arguments in favour of $\overline{MS}$ regularization scheme are that from one side the calculated up to now higher order radiative corrections in $\overline{MS}$ scheme are smaller than in MS scheme and from the other side that the calculations in MS scheme are simpler than in momentum regularization scheme. Several years ago Stevenson had suggested another appoach to the problem, which he called the Principle of Minimal Sensitivity (PMS) [122]. According to Stevenson the renormalization scheme must be chosen for each process separately from the requirement, that physical quantity, which is independent on renormalization scheme when all terms of perturbation theory are summed, must have the minimal dependence on the scheme in any order of perturbation theory. Recently a scheme independent method of regularization is proposed [123-125], where the coefficients of perturbation serie for a given matrix element are expressed in terms of the value of the same matrix element. In the presented at the Conference report [126] this method was generalized to the case of massive quarks. It was shown also, that the results of this approach and PMS scheme are very close to one another.

But what says experiment about the numerical values α_s and Λ_{QCD}? Comparing with the Lepage report at the Cornell symposium [105] where this problem was discussed in detail, little new happened. It became even more clear that finding of α_s from the measurements of jets or multihadron events in e^+e^- annihilation strongly depends on fragmentation models [106,107]. As was mentioned, because of large nonperturbative corrections hadronic charmonium decays cannot be used for determining α_s. The data on upsilonium hadronic decays give [65] $\alpha_s(m_b) = 0.165 \pm 0.005$ and $\Lambda_{\overline{MS}} = 118^{+16}_{-15}$ MeV.

The spead of experimental values of $\Lambda_{\overline{MS}}$ found in deep-inelastic lepton-nucleon scattering is still large $\Lambda_{\overline{MS}}$=50-300 MeV. But probably they do not contradict the value stated above (see I.A. Savin plenary talk).

III. Conclusion

In the last few years we became witnesses of an impressive progress in QCD: from an abstract theoretical scheme it transforms into a working tool. A lot of beautiful physics has been understood and explained but it remains to understand even more.

QCD is nowadays close to solving very important problems of strong interaction physics - to foundate and to find applicability limits of various model approaches. One of these approaches - soft pion theory follows directly from QCD and, as I have already told, QCD allows one to determine phenomenological parameters figuring in it. So that in this case the problem is close to its solution. In a number of cases we are able to verify and specify with QCD the vector dominance model.

What has not been done and what, in my opinion, is extremely important - is establishment of relationship between QCD and Regge approach in constructive way - QCD calculation of Regge trajectories, residues, finding of applicability limits of Regge theory at large t etc.

Another important problem is derivation of constituent quark model from QCD. At present, where such tests were managed to be done, QCD confirmed quark model conclusions with more or less accuracy. However, QCD and quark model proceed from quite different starting points: massless quarks, chirality conservation, and, generally speaking, strong spin dependence (QCD) and massive quarks and SU(6) symmetry, i.e. the absence of spin dependence (quark model). By virtue of this, coincidence of QCD and quark model results is rather unnatural: strongly spin-dependent zero approximation terms in QCD sum rules are compensated by large power corrections. It would be desirable to have here a better understanding or, at least, to find effects where predictions of QCD and of constituent quark model strongly diverge.

Unfortunately, there are yet no ideas how to describe radial expectations in QCD - the constituent model has here a great advantage over QCD. This is also connected with the fact that the glueball problem is yet unclear. To solve this problem one must be in a position to determine the high-excited quark and gluon state mixing.

In perturbative QCD the Sturm und Drang period beginning after discovery of asymptotic freedom gave way in the last few years to a period of a more slow development. It was understood that previously disregarded terms of higher order in α_s and higher power corrections were important in many cases. It became clear also that definite theoretical predictions of observable phenomena in most cases cannot be achieved without better understanding of the mechanism of quark and gluon transition into hadrons. One can hope, however, that if the relationship between QCD and Regge approach will be established, then it will throw a new light on this problem.

Last but not least what is needed is a satisfactory model of vacuum which could permit a reliable calculation of various vacuum expectation values and put a firm basis under to-day's nonperturbative QCD calculations.

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Discussion:

C.A. Dominguez: I would like to comment on your predictions for the mass spectrum. The values of m_s and $\langle\bar{s}s\rangle/\langle\bar{u}u\rangle$ that you need are hard to reconcile with those determined from the pseudoscalar sum rules. I wish to stress that by improving the continuum model including pseudoscalar radial excitations and using power corrections from up to dimension six it is indeed possible to meet your criterion that predictions be bounded from below as well as from above. The results confirm the early Marseille bounds on m_s as well as Narison's estimate of $\langle\bar{s}s\rangle/\langle\bar{u}u\rangle$.

B.L. Ioffe: As I have said in my talk the QCD sum rule approach does not work in pseudoscalar channel. The easiest way to see this is to consider nondiagonal polarization operator $\langle 0|T\{\bar{u}\gamma_5 u, \bar{s}\gamma_5 s\}|0\rangle$ and compare it with $\langle 0|T\{j_8, j_8\}|0\rangle$, $j_8=(\bar{u}\gamma_5 u+d\gamma_5 d-2\bar{s}\gamma_5 s)/\sqrt{6}$. The contribution of physical states to both are of the same order (η-meson), but in QCD sum rule approach the first is much less, than the second. Novikov, Shifman, Vainstein and Zakharov had shown by direct calculation that the discrepancy between left and right hand sides of QCD sum rules persist also in the case of pseudoscalar isovector (pion) current. Therefore I think that the determination of m_s and $\langle\bar{s}s\rangle/\langle\bar{u}u\rangle$ by consideration the pseudoscalar sum rules is not reliable.

S. Narison: My comments concern the result for $\langle\bar{s}s\rangle/\langle\bar{u}u\rangle-1 \simeq -0.2^{+0.05}_{-0.1}$ and other results from the baryon sum rules: these results should be affected by the large radiative connections, the choice of the baryon operators and the higher excitations as discussed recently by the Heidelberg group in the (u,d) channel. Actually, your choice of operators for the description of the baryons might not be the optimal one. So, I think that your error bars for $\langle\bar{s}s\rangle/\langle\bar{u}u\rangle$ are insignificant and your results for $\langle\bar{s}s\rangle/\langle\bar{u}u\rangle$ might still be consistent with the ones from the scalar and pseudoscalar meson sum rules.

B.L. Ioffe: The radiative corrections mostly affects the residues, in lesser extent the proton mass and only slightly influence the mass differences from which $\langle\bar{s}s\rangle/\langle\bar{u}u\rangle$ was determined. Ofcourse, the calculations of α_s corrections to baryon sum rules would be very important, but up to now no definite answers is known: there is a discrepancy between results of Heidelberg and INR groups I will not to repeat here the arguments for my choice of baryon current, they are published in my paper in Zs. f. Physik, C. Now this choice is supported by the whole set of results obtained with this current. The recent critizism of my choice of baryon current by Heidelberg group is not convincing since it is based on the requirement of very fine concidence of left and right hand side of sum rules, which is nonsense in sum rule approach.

V.A. Rubakov: I would like to mention that within the last few years the alternative variants of QCD sum rules were developed. Namely, the Finite Energy Sum Rules were considered by Kataev, Krasnikov and Pivovarov from our Institute and it turned out that FESR are particularly convenient for calculating the properties of radial excitations of mesons. Another development is the invention of the so called generalized sum rules by Chetyrkin, Krasulin and Matveev (also from INR). This generalization is necessary for evaluating quantities like the pion form-factor at small momentum transfer (which has been calculated by this group and Radiushkin and Nesterenko).

B.L. Ioffe: The Finite Energy Sum Rule (FESR) approach to calculation of hadronic masses differs from QCD sum rule one at only one point: in the former a cut off of dispersion integrals is used instead of consideration of Borel (Laplace) transformed equations in the latter. In all other aspects the FESR approach to problems in view is simple a repetition of what was done by QCD sum rule method. The dispersion integrals, which appears here are divergent (and rather strongly in some cases). The cut-off of such integrals cannot be justified and therefore the accuracy of the obtained results is not under control. (In some cases there is a strong disagreement with QCD sum rule results: by a factor ~ 1.5 in masses and 10 in residues, even if the cut-off parameter is chosen by a best fit method.)