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QCD SPECTROSCOPY AND COUPLINGS WITH SUM RULES: AN OVERVIEW*)

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ABSTRACT

An overview is presented of the results for masses and couplings of resonances obtained in recent years with QCD sum rules.

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In this paper we will review the status of the QCD sum rule approach to hadronic physics. Details of the calculations can be found in the original papers referred to below, in review papers already in the literature or in a Physics Reports which is at present in preparation.

We will discuss many topics and will try to give a coherent view of the present situation. We will answer criticism and discuss outstanding questions and problems.

THE SVZ EXPANSION

Several years ago Shifman, Vainshtein, and Zakharov[1)] showed how to generalize the short-distance expansion for theories with a non-trivial vacuum. Although the method is very appealing physically, it is not clear that the prescription suggested by the ITEP people is indeed correct. Phenomenologically, as can be seen in the accompanying tables, the success is impressive. Theoretically there is considerable resistance to this approach. There are some situations of solvable models in which the prescription works, e.g. expansion in an external instanton field[2)]. Also, contrary to initial claims[3)] the solution of ϕ^4 in the perturbative vacuum in the presence of condensates agrees with the SVZ prescription[4)].(the answer is the same before and after shifting). There are studies in other models by David[5)] which raise other difficulties. Essentially, it is claimed that condensates which are not protected by a symmetry are necessarily ill-defined. We believe that the question is not settled.

SPECTROSCOPY

The calculation of masses works schematically as follows. In the Euclidean region one calculates the polarization function for a current with a well defined set of quantum numbers that single out a particular partial wave. For example, the current $\bar{u}\gamma_\mu\bar{u} - \bar{d}\gamma_\mu d$ for light quark vector states like the ρ meson $(I = 1,\ J^{PC} = 1^{--})$, while the current $\bar{c}\gamma_\mu c$ would select vector states made of charmed quarks. The momentum Q^2 flowing through the current is chosen in the asymptotically free region. While in the examples given above the current is physical and measured in e^+e^- annihilation, other currents will be unphysical.

Applying the SVZ expansion to the polarization function, one obtains an expression of the form

$$T_{\mu\nu\ldots}\Pi^j(Q^2) = i\int d^4x\, e^{iqx} \langle 0|T(j_1(x)j_2(0)|0\rangle = \sum_i a_i \langle 0|O_i|0\rangle, \quad Q^2 = -q^2 \qquad (1)$$

$\Pi^j(Q^2)$ is a scalar function of Q^2; $T_{\mu\nu}\dots$ a tensor depending on the current in question, the O_i are local operators constructed of quark and gluon fields, and the a_i the corresponding Wilson coefficients. Equation (1) represents the theoretical side of the sum rules and is calculated in terms of the fundamental parameters of QCD (α_s, quark masses) and of condensates (vacuum matrix elements of the operators O_i). The dynamics of QCD is buried in the Wilson coefficients a_i. These can be calculated in perturbation theory and their magnitude determines the validity region of the expression, which is only applicable if the neglected terms are indeed small. The first operator in (1) is the identity operator. Its coefficient a_o contains the ordinary perturbative contributions and in most cases is calculated to first order in α_s. Therefore, one has to make sure that higher corrections are small. For dimensional reasons the coefficients of the higher dimensional operators in (1) fall off by corresponding powers of Q^2. These power corrections measure the breakdown of asymptotic freedom and grow fast with distance. The theory makes sense if the power corrections are big enough for quarks to resonate and small enough so that contributions from higher dimensional operators can be neglected. The balancing of the various corrections is a tricky path and people keep falling into the precipice[6]. In most cases the errors can be calculated and although the method cannot solve all problems its range of validity can always be tested.

The phenomenological side that matches the theoretical expression has the following simple form in the physical region

$$\sum_i \frac{g_i^2}{m_i^2 - Q^2} + \text{continuum} \tag{2}$$

where the (g_i, m_i) are the parameters of the resonances. The continuum in (2) contains the parameter s_o which is the threshold at which the smooth background starts.

After analytic continuation to the Euclidean region via the moment or Borel method[1] the two sides hopefully overlap. Using one or the other method depends on whether the problem has a natural scale (like a heavy quark) or not. Nothing basic is involved in the choice of method. For massless quarks it is just inconvenient to use a reference momentum as a scale. The crucial question is whether one is gaining physical insight and information, or is just trading masses and couplings for new parameters like condensates and thresholds.

The spectroscopy can be divided in a number of sections:

1. light quark mesons with L = 0[1)]
2. light quark mesons with L = 1[7),8)]
3. charmonium[1),9)]
4. bottonium[9),10)]
5. mesons made of light and heavy quarks[11),12)]
6. light quark baryons[13-17)]
7. baryons with one heavy quark[18)]
8. QCD non-quark model states[19-20)]

We will briefly give a status report for each case with emphasis on parameter dependences.

1. Light quark mesons with L = 0[1)]

The operator expansion series (1) for the polarization operator can be cut off at dimension d = 6 for the operators O_i. Consequently, the operators which give important contributions to these mesons are (apart from the identity operator)

$$\langle 0|\frac{\alpha_s}{\pi} G^a_{\mu\nu} G^a_{\mu\nu}|0\rangle \quad \text{the gluon condensate,} \quad d = 4$$

$$\langle 0|m\bar{q}q|0\rangle \quad \text{the quark condensate,} \quad d = 4 \tag{3}$$

$$\langle 0|\bar{q}\Gamma_1 q\bar{q}\Gamma_2 q|0\rangle \quad \text{four-fermion condensate,} \quad d = 6$$

Assuming dominance of the vacuum intermediate state[1)] the last operator in (3) can be expressed in terms of $\langle 0|\bar{q}\, q|0\rangle$. This is claimed to be accurate within about 10%[21)]. However, due to uncertainties in the light quark mass values, $\langle 0|\bar{q}\, q|0\rangle$ is only known up to a factor of two (from PCAC).

Consider as an example the resulting expression for the polarization function of the ρ-meson current (after Borel transforming)

$$\int e^{-s/M^2} \mathrm{Im}\Pi(s)\, ds = \frac{1}{8\pi^2} M^2 \Big[1 + \frac{\alpha_s(M)}{\pi} + \frac{8\pi^2}{M^4}\langle 0|m\bar{q}q|0\rangle + \frac{\pi^2}{3M^4}\langle 0|\frac{\alpha_s}{\pi} G^a_{\mu\nu} G^a_{\mu\nu}|0\rangle - \frac{448}{81}\cdot\frac{\pi^3\alpha_s}{M^6}\langle 0|\bar{q}q|0\rangle^2\Big] \tag{4}$$

A second sum rule can be derived by differentiating with respect to M^2. An expression for the ρ-meson mass can be obtained by substituting

$$\mathrm{Im}\Pi(s) = \frac{\pi m_\rho^2}{g_\rho^2}\,\delta(s - m_\rho^2) + \frac{1}{8\pi}\left(1 + \frac{\alpha_s}{\pi}\right)\theta(s - s_0) \tag{5}$$

into (4) and into the derivative sum rule, transferring the continuum contribution to the right-hand side and taking the ratio

$$m_\rho^2 = M^2 \frac{(1+\frac{\alpha_s}{\pi})\left[1-(1+\frac{s_0}{M^2})e^{-s_0/M^2}\right] - \frac{\pi^2}{3M^4}\langle\frac{\alpha_s}{\pi}G^2\rangle + \frac{896}{81}\cdot\frac{\pi^3\alpha_s}{M^6}\langle\bar{q}q\rangle^2}{(1+\frac{\alpha_s}{\pi})\left[1-e^{-s_0/M^2}\right] + \frac{\pi^2}{3M^4}\langle\frac{\alpha_s}{\pi}G^2\rangle - \frac{448}{81}\cdot\frac{\pi^3\alpha_s}{M^6}\langle\bar{q}q\rangle^2} \qquad (6)$$

where we have neglected the contribution from $m\bar{q}q$ which is small compared to the gluon condensate. It can be seen from (6) that the perturbative corrections $(1+\frac{\alpha_s}{\pi})$ do not play a role as they can be divided out at the expense of introducing a 20% error in the values for the condensates.

In (6) the mass m_ρ^2 is controlled by the gluon and quark condensates on an equal footing, and although the ρ' is far away at 1.6 GeV it still contributes about 30% to the sum rule at $M^2 \simeq m_\rho^2$. This contribution is controlled by s_o. Due to the uncertainties in s_o and $\langle\bar{q}q\rangle$ it is not possible to determine the gluon condensate $\langle\frac{\alpha_s}{\pi}G^2\rangle$ from (6) better than within a factor of two. The sum rule can also be analyzed without ρ dominance using the e^+e^- annihilation data. This analysis[22)] confirms the values of the matrix elements to this accuracy. One can use the sum rule (6) to determine the parameters of the ρ meson. For $s_o = 1.5\ \mathrm{GeV}^2$ which is suggested by the data, one gets

$$m_\rho = 770\ \mathrm{MeV} \pm 10\%$$

$$g_\rho^2/4\pi \simeq 2.5$$

Similar sum rules can be derived for all other vector mesons. Their masses and couplings can be calculated with the same accuracy. The pattern of SU(3) breaking comes out correctly. The mass of the strange quark is not accurately fixed by these sum rules. In Table I the results have been summarized.

The pseudoscalars are too light for the method to be able to determine their masses. Moreover, direct instanton contributions to the pseudoscalar sum rules not included in the matrix elements $\langle O_i\rangle$ of Eq. (1) have to be accounted for[23)]. Using axial vector sum rules at $M^2 \simeq m_\rho^2$ where on one hand the power corrections are small and on the other hand the pseudoscalars saturate the sum rule, one can obtain the couplings in good agreement with experiment. Also, the consistency of the axial sum rules for the A_1 demands the pion pole at zero.

2. The L = 1 states[7),8)]

The L = 1 states with massless quarks give about 15 new predictions with the same operators and parameters. Although these states are more sensitive to the continuum, reasonable variations in s_0 do not change the results by more than 10%. As can be seen in Ref. 7), the power corrections are always important and a gluon condensate which is a factor two bigger than the standard value would completely ruin the nice agreement for spin 2 mesons. This is indirect evidence for the standard value. Also, these states require m_s to be small (about 100 MeV). The results are collected in Table II.

3. Charmonium[1),9)]

This system isolates the gluon condensate. Moreover, for kinematical reasons everything conspires to give maximal accuracy. This is achieved by allowing the reference momentum Q^2 to vary in the region where asymptotic freedom is valid. Besides getting a better stability region the results are stable for a wide range of Q^2 [Ref. 9] which is strong support for the theory. Using moment sum rules one can make sure that the continuum contribution is unimportant, which can be verified from the measured cross-section in the vector current case. The mass of the charmed quark is accurately fixed and the gluon condensate is pinned down with about 30% accuracy and at the same value as in light quark spectroscopy. The masses come out naturally and accurately. The results are collected in Table III. Corrections due to higher operators are very small for dimension six operators and although large for dimension eight operators at $Q^2 = 0$,[24)] their contribution is expected to go down when Q^2 is shifted away from zero where the stability is best[9)]. This calculation has to be performed and if our expectation is confirmed, this will corroborate the results in a strong way.

4. Bottonium[9),10)]

This system is unfortunately on the limit of being Coulombic and therefore very different from charmonium. Since the gluon condensate coefficient is down by the ratio of the quark masses to the fourth power, it is clear that low moments are dominated by perturbative corrections. To feel the resonances one has to go to large distances, but then higher order perturbative corrections are also not negligible. It is only when the quark mass is about 20 GeV that the system becomes tractable again[25),26)]. A few things can nevertheless be said (see Table III) and with further assumptions one can make predictions (but at a risk)[10)].

5. Heavy-light quark bound states[11),12)]

Unfortunately, moment sum rules cannot be used to obtain the masses of open charm states. As stated in the Introduction, conditions on the size of the various contributions must be met for the polarization operator formula to be valid. Using the reference momentum in a range where asymptotic freedom is valid one can tune these contributions. At the same time, however, one increases the contribution from the continuum to such an extent that the lowest lying resonance cannot be disentangled from the background.

For open bottom the situation is better and the results show quite interesting physics. The quark mass is fixed from the sum rules for bottonium and the only parameter is the continuum threshold s_o. A stable set of states which depend on s_o is obtained[11)]. The most interesting result is that the P-wave mesons are higher than in potential models because of the appearance of terms of the form $m_Q \langle\bar{q}q\rangle$ where m_Q is the heavy quark mass.

Recently[12)] Borel transformed sum rules for open charm and beauty have been analyzed in order to obtain the couplings f_D and f_B which within errors are equal to f_π. In this case the experimental mass values were used as input. Although there is a large uncertainty in the values obtained, these results and those of Ref. 11) rule out the large value for f_D that has been conjectured in the charm case to explain the $D^+ - D^o$ lifetime[33)].

6. Baryons

Light quark baryons are nicely understood in this language, although it is very difficult to obtain high accuracy. For the $J = 1/2$ baryons there are two structure functions $A(q^2)$ and $B(q^2)$, which have odd and even numbers of dimensions respectively. Consequently, the expansion for the odd dimensional function $A(q^2)$ is dominated by $\langle 0|\bar{q}q|0\rangle$ which appears without the mass of the light quark. Approximating one sum rule (from $A(q^2)$) by the quark condensate contribution only, the other (from $B(q^2)$) by the bare loop, and taking the ratio one finds the celebrated formula of Joffe[13)]

$$M_N^3 = -2(2\pi)^2 \langle 0|\bar{q}q|0\rangle \tag{7}$$

Using the PCAC value of $\langle 0|\bar{q}q|0\rangle$ one obtains $M_N \cong 1$ GeV (very sensitive to the value of $\langle 0|\bar{q}q|0\rangle$). This result can be improved by including higher dimensional operators and the continuum contribution but results at the 10% accuracy level or better, which are necessary to reproduce the splittings within the octet must wait some further calculations and a more accurate determination of $\langle 0|\bar{q}q|0\rangle$ which has a rather large uncertainty[27].

Nevertheless, there is considerable qualitative agreement with experiment and with theoretical prejudice. The nucleon mass vanishes with the condensate. The lowest lying states fit into a positive parity 56 representation while negative parity states come out higher[16]. It turns out that the octet baryons provide an effective way to establish the value of

$$\gamma = -1 + \langle 0|\bar{s}s|0\rangle / \langle 0|\bar{u}u|0\rangle$$

irrespective of the values of the other parameters[17]. Provided the strange quark mass is between 100 and 150 MeV, one finds $\gamma = -0.18 \pm 0.06$. The sign does not agree with extrapolations from chiral perturbation theory which, however, is not surprising since the quark mass is of order Λ_{QCD}[23].

The decuplet masses are marginally within reach of the method since the non-perturbative contributions are significantly larger in this case. Qualitatively, however, the agreement is satisfactory.

In general the baryon spectrum is dense which implies that s_o is low and that the baryons saturate less than 50% of the sum rule. It will take more detailed calculations to reach 10% or better accuracy. In particular, the perturbative gluon exchange contributions have not yet been taken into account. Finally, the dimension five operator $\bar{q}\sigma_{\mu\nu}q\, G_{\mu\nu}$ gives an important contribution to the octet-decuplet splitting. Its matrix element is not well known and no hard prediction is possible here.

Finally, the coupling λ_N of the proton to the current which measures the strength of the proton transition into three quarks has been used[34] to calculate the proton lifetime in the SU(5) model of grand unification.

7. Heavy-light baryons

The situation here is quite open. There are some calculations in the literature[18] based on rough approximations. We do not believe that these test the theory, but it is reassuring that with the known mass of the charmed quark the charmed baryon appears at the expected mass.

8. Exotic states

It is rather distressing that the gluonic degrees of freedom of QCD are not seen. Some calculations of glueballs and of matter with glue have been performed with QCD sum rules.

The ITEP group[23)] has made a semiquantitative sum rule analysis for glueballs and concludes that they should be relatively high in mass and that the spin 2 glueball lies lowest. Experimentally the situation is still confused, but it is important to emphasize that bag models give different results.

For mesons with glue there exist two calculations in the literature[19),20)]. They disagree and therefore must be checked. One of them[20)] makes unnecessary assumptions on subtractions but claims that the Wilson coefficients in Ref. 19) are not correct. Both predict a low mass for the exotic $\bar{q}qG$ state with $J^{PC} = 1^{-+}$ which is worrying. We believe that experiments searching for these states are as interesting as glueballs and should be pursued.

Similar calculations for baryons with glue are in progress but experimentally there seems to be little room for more states below 2 GeV as predicted by QCD like models and bags[28)]. If the sum rules give similar results, it is important to understand the reasons.

9. Three-point functions

The use of these methods for three-point functions was pioneered by the ITEP group in 1978[29)] and developed to new realms since.

One application is to electromagnetic processes like $\eta_c \to 2\gamma$[30)] and $\psi' \to \eta_c\gamma$[31)]. The anomaly diagram is used here in a different domain. By making a moment expansion (or a double moment expansion for $\psi' \to \eta_c\gamma$) one looks for a region in which the relevant states dominate. Combining the results with couplings determined from two-point function analyses one can extract the transition rates. Unfortunately, the perturbative piece in the case of $\psi' \to \eta_c\gamma$ cannot be calculated. However, it is quite remarkable that the non-perturbative gluon condensate contribution substantially reduces these rates, bringing them in better agreement with experiment.

Another interesting application is the calculation of strong interaction coupling constants like $g_{\pi NN}$ and $g_{\omega\rho\pi}$[32)]. These calculations have been performed by applying the Wilson expansion to the product of three currents and can be checked by considering two-point functions sandwiched between the vacuum and a pion state. The procedure followed for three-point functions is to isolate the pion pole term and analytically continue the residue. This procedure is consistent with chiral invariance and the relations obtained are equivalent to the Goldberger-Treiman relation or current algebra results. We obtain for example:

$$g_{\pi NN} \simeq 2(2\pi)^2 \sqrt{2}\, \frac{f_\pi}{M_N}$$
$$g_{\omega\rho\pi} \simeq \sqrt{2}\,(2\pi)^2\, \frac{f_\pi}{m_\rho^2} \qquad (8)$$

The numerical values of these and other couplings are tabulated in Table IV.

Finally we have to mention the applications of the three-point function technique to the calculation of form factors[35),36)] and an alternative method to calculate coupling constants[37)].

Table I

Light quark mesons with L = 0 (Ref. 1)

State	J^{PC}	Mass exp	Mass theor	Coupling exp	Coupling theor	Remarks
π	0^{--}	140	-	f_π=133	125	masses too low direct instantons f_π, f_K well computed.
K		495	-	$f_K \simeq f_\pi$		
η		550	-			
η'		920	-			
				$g^2/4\pi$		masses and couplings calculated with $\sim$ 10% acc. ρ-ω interference has also been obtained. Relevant parameters are $\langle G^2\rangle$, $\langle\bar{q}q\rangle$, α_s, s_o, $m_q \equiv 0$, I = 0,1 degeneracy, SU(3) breaking o.k., m_s not well fixed.
ω	1^{--}	780	770	2.4	2.5	
ρ		770	770	2.4	2.5	
K^*		890	930	1.39	1.4	
ϕ		1020	1070	12.0	14.	

Table II

Light quark mesons with $L = 1$ (Ref. 7, 8)

J^{PC}	state	mass exp	mass theor	coupling exp	coupling theor	Remarks
2^{++}	A_2 f K^{**} f'	1320 1270 1430 1520	1320 1320 - 1520	couplings not directly useful; g_f=0.04 calculated with further assumption, agrees with exp.		I = 1,0 degeneracy. m_s low ∿ 120 MeV. Large $1/M^4$ term gives bound on $\langle G^2 \rangle$.
1^{++}	A_1 D E Q_1 Q_2	1200 1285 1420 1270 1414	1150 1290 1460 - -	$\frac{4\pi}{f^2_{A_1}}$ =.15	.16	two sum rules for A_1; D meson requires m_s ∿ ∿ 100 MeV; for D only one sum rule, since divergence of axial current has U(1) problem in this channel.
0^{++}	δ S^* ε	980 980 1300	1010 1010 1350			δ, S^* assumed to be pure $\bar{q}q$, no instanton contributions included.
1^{+-}	B	1240	?			no calculation possible; power corrections vanish at one-loop level
2^{-}	A_3	1680	1630			

Table III

Heavy quark mesons with L = 0 and L = 1

J^{PC}	State	mass (GeV) exp.	theor	Coupling	Remarks
Charmonium (Ref. 1,9)					
1^{--}	J/ψ	3.10	3.10±0.01	Only for J/ψ: exp: $\Gamma_{e^+e^-}$ = 4.7 ± .6 keV theor: $\Gamma_{e^+e^-} \simeq$ 5.34 keV	$m_c(p^2=-m_c^2)=1.28$ GeV very accurate indep. of s_o. Gluon condensate same as in light quark case $\langle\frac{\alpha_s}{\pi} G^2\rangle \simeq (330\ \mathrm{MeV})^4 \pm 30\%$. Same parameters fit P-waves.
0^{--}	η_c	2.98	3.01±0.02		
0^{++}	χ_0	3.42	3.40±0.01		
1^{++}	χ_1	3.51	3.50±0.01		
2^{++}	χ_2	3.56	3.56±0.01		
1^{+-}		?	3.51±0.01		
Bottonium (Ref. 9)					
1^{--}	Υ	9.46	$m_\Upsilon - m_{\eta_b} \sim$ ~ 60 MeV		$m_b(p^2=-m_b^2) \simeq 4.26$ GeV moment method fails, no single resonance saturation
0^{--}	η_B	?			
Open bottom (Ref. 11,12)					
0^{-+}	$(\bar{u}b)$	5.27	5.31	$f_P \simeq$ 140-180 MeV	continuum very important; splittings cannot be resolved.
1^{--}			5.38	$g_V^2/4\pi \simeq 24$	
0^{++}			6.13	$f_S \simeq$ 380 MeV	S-P splitting large because of $m_Q\langle\bar{q}q\rangle$.
1^{++}			6.17	$g_A^2/4\pi \simeq 13$	
0^{-+}	$(\bar{s}b)$		5.42	$f_P \simeq$ 320 MeV	
1^{--}			5.46	$g_V^2/4\pi \simeq 17$	
0^{++}			6.29	$f_S \simeq$ 400 MeV	
1^{++}			6.34	$g_A^2/4\pi \simeq 12$	

Table IV

Couplings $g^2/4\pi$ of Goldstone bosons to hadrons (Ref. 32)

	exp[38)]	theory	
πNN	14.5	12.5	all within 20% of
$\pi\Sigma\Sigma$	13±2	10	exp value
$\eta_8 NN$	~4.5	6.4	$\alpha = D/(F+D) = 7/12$
$K\Sigma N$	~1	1	(compare $\alpha(SU/6)) = 2/3$)
$\pi N\Delta$	~15 GeV^{-2}	18 GeV^{-2}	
$\omega\rho\pi$	~16 GeV^{-1}	13 GeV^{-1}	

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