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THE PION WAVE FUNCTION AND QCD SUM RULES WITH NONLOCAL CONDENSATES [1]

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Abstract

The QCD sum rule calculation of the pion wave function by Chernyak and Zhitnitsky is implicitly assuming that the correlation length of vacuum fluctuations is large compared to the typical hadronic scale $\sim 1/m_\rho$, so that one can substitute the original nonlocal objects like $\langle\bar{q}(0)q(z)\rangle$ by constant $\langle\bar{q}(0)q(0)\rangle$-type values. We outline a formalism enabling one to work directly with the nonlocal condensates, and construct a modified sum rule for the moments $\langle\xi^N\rangle$ of the pion wave function. The results are rather sensitive to the value of the parameter $\lambda_q^2 = \langle\bar{q}D^2q\rangle/\langle\bar{q}q\rangle$ specifying the average virtuality of the vacuum quarks. Varying it from the most popular value $\lambda_q^2 = 0.4\,GeV^2$ up to the value $\lambda_q^2 = 1.2\,GeV^2$ suggested by the instanton liquid model, we obtain $\langle\xi^2\rangle = 0.25 - 0.20$, to be compared to the CZ value $\langle\xi^2\rangle = 0.43$ obtained with $\lambda_q^2 = 0$.

1.Introductory remarks. The standard trick incorporated in all the approaches based on the asymptotic freedom of QCD and factorization is the introduction of some phenomenological functions and/or numbers accumulating necessary information about nonperturbative long-distance dynamics of the theory. The most important examples are:

- parton distribution functions $f_{p/H}(x)$ used in the perturbative QCD approaches to hard inclusive processes [1],
- hadronic wave functions $\varphi_\pi(x)$, $\varphi_N(x_1, x_2, x_3)$, *etc.*, which naturally emerge in the asymptotic QCD analyses of hard exclusive processes [2, 3, 4, 5, 6]
- quark and gluon condensates $\langle\bar{q}(0)q(0)\rangle$, $\langle G(0)G(0)\rangle$, the basic parameters of the QCD sum rule approach [7], describing the nonperturbative nature of the QCD vacuum.

The hope is that in some future approach they all will be calculated from the first principles of QCD without any model and/or *ad hoc* assumptions. A less ambitious program is to calculate the hadronic functions $f(x), \varphi(\{x\})$ using the QCD sum rules [7], with only the condensate values treated as input parameters.

While the parton distribution functions can be extracted rather reliably from experimental data, the situation with the hadronic wave functions is much more complicated. Normally, they appear only in an integrated form. Furthermore, the very applicability of the perturbative QCD formulas at accessible energies is questionable [8, 9]. In this situation, the QCD sum rule approach and lattice calculations are the only reliable way to get information about the form of the hadronic wave functions. In particular, the most popular set of hadronic wave functions [10], due to Chernyak, A.Zhitnitsky and I.Zhitnitsky (CZ), was produced with the help of QCD sum rules.

One should remember, however, that the operator product expansion, the starting point of any QCD sum rule analysis, has different forms depending on the situation. Presence of a large (or small) extra parameter might essentially modify the expansion. The most well studied example is the modification of the OPE for the form factors at small momentum transfer q [11, 12]. In that case a simple-minded extrapolation from the region of moderately large q is completely unjustified: one cannot reproduce in that way even the normalization conditions like $F_\pi(0) = 1$. Our goal in the present paper is to show that calculating the $N \geq 2$ moments of the pion wave function one faces another situation requiring a modification of the underlying expansion. We construct a modified sum rule and show that, for a standard choice of the condensate values, it produces the pion wave function that strongly differs from the CZ form.

2. Pion wave function and QCD sum rules: criticism of the CZ-approach. The first application of the QCD sum rules to the pion wave function $\varphi_\pi(x)$ was the calculation of its zero moment, *i.e.*, the pion decay constant f_π , in the pioneering SVZ paper [7]. It was calculated there within 5% accuracy. This success inspired Chernyak and A.Zhitnitsky [13] to calculate the whole pion WF by reconstructing it from the next moments $\langle\xi^N\rangle$ (where $\xi = 2x - 1$). They extracted $\langle\xi^2\rangle$ and $\langle\xi^4\rangle$ from the relevant SR precisely in the same way

f_π value. However, the nonperturbative terms in their sum rule

$$f_\pi^2\langle\xi^N\rangle = \frac{M^2}{4\pi^2}\frac{3}{(N+1)(N+3)}(1-e^{-s_0/M^2}) + \frac{\alpha_s\langle GG\rangle}{12\pi M^2} + \frac{16}{81}\frac{\pi\alpha_s\langle\bar{q}q\rangle^2}{M^4}(11+4N) \qquad (1)$$

completely different N-dependence compared to the perturbative one and, *a priori*, it lear whether a straightforward use of the $N=0$ technology can be justified for higher e scale determining the magnitude of all the hadronic parameters including s_0 (the uum threshold" [7]) is eventually settled by the ratios of the condensate contributions perturbative term. If the condensate contributions in the CZ sum rule (1) would have ne N-behavior as the perturbative term, then the N-dependence of $\langle\xi^N\rangle$ would be ined by the overall factor $3/(N+1)(N+3)$ and the resulting wave function $\varphi(x)$ coincide with the "asymptotic" form [4, 6]

$$\varphi_\pi^{as}(x) = 6f_\pi x(1-x). \qquad (2)$$

wever, the ratios of the $\langle\bar{q}q\rangle$ and $\langle GG\rangle$-corrections to the perturbative term in eq. (1) wing functions of N. In particular, in the $\langle\bar{q}q\rangle$ case, the above mentioned ratio for is by factor 95/11 larger than that in the $N=0$ case. For $N=4$ the enhancement equals 315/11 . As a result, the effective vacuum scales of $(mass)^2$ dimension are by $(95/11)^{1/3}\approx 2.1$ and $(315/11)^{1/3}\approx 3.1$ larger than that for the $N=0$ case. Ap- ately the same factors ($5^{1/2}\approx 2.2$ and $(35/3)^{1/2}\approx 3.4$) one obtains also for the gluon sate term. Hence, the parameters $s_0^{(N)}$ and the combinations $f_\pi^2\langle\xi^N\rangle$ straightforwardly ed from the SR (1) must be larger than the "asymptotic" values $s_0^{N=0}\approx 0.75\,GeV^2$ $\langle\xi^N\rangle^{as} = 3f_\pi^2/(N+1)(N+3)$ just by the factors 2 (for $N=2$) and 3 (for $N=4$). are just the results obtained in Ref.[13].

better understand the structure of the relevant power series it is instructive to rewrite for the pion wave function $\varphi_\pi(x)$ itself [14]:

$$\begin{aligned} f_\pi^2\varphi_\pi(x) &= \frac{M^2}{4\pi^2}(1-e^{-s_0/M^2})\varphi_\pi^{as}(x) + \frac{\alpha_s\langle GG\rangle}{24\pi M^2}[\delta(x)+\delta(1-x)] \\ &\quad + \frac{8}{81}\frac{\pi\alpha_s\langle\bar{q}q\rangle^2}{M^4}\{11[\delta(x)+\delta(1-x)] + 2[\delta'(x)+\delta'(1-x)]\}. \qquad (3) \end{aligned}$$

e $O(1)$ and $O(N)$ terms in eq. (1) correspond to the $\delta(x)$ and $\delta'(x)$-terms in eq.(3). In a, presence of the $\delta(x)$-functions in eq.(3) is evidently indicating that the vacuum fields ated as carrying zero fraction of the pion momentum. This can be easily understood erving that the operator product expansion (underlying eqs.(1),(3)) is, in fact, a power expansion over small momenta k of vacuum quarks and gluons. Retaining only the nd $\langle GG\rangle$-terms (like in eqs.(1),(3)) is just equivalent to the assumption that k is not small but exactly equals zero.

wever, it is much more reasonable to expect that the vacuum quanta have a smooth ution with a finite width μ. In configuration space, this means that vacuum fluctua- ave a finite correlation length of the order of $1/\mu$, so that the two-point condensates

The expansion of the condensate $M(z^2)$ over the local condensates corresponds to that of the distribution function $f_S(\nu)$ over the $\delta^{(n)}$-functions:

$$f_S(\nu) = \delta(\nu) - L_S\delta'(\nu) + \dots , \tag{7}$$

with L_S fixed just by the average virtuality of the vacuum quarks (eq.(4)): $L_S = \lambda_q^2/2$.

There is another (vector) bilocal condensate $M_\mu \equiv \langle\bar{q}(0)\gamma_\mu q(z)\rangle$, containing a γ-matrix:

$$\langle\bar{q}(0)\gamma_\mu q(z)\rangle = iz_\mu A\int_0^\infty e^{\nu z^2/4} f_V(\nu)\,d\nu \tag{8}$$

where $A = \frac{2}{81}\pi\alpha_s\langle\bar{q}q\rangle^2$. The zeroth moment of $f_V(\nu)$ is zero in the limit of massless quarks, and that is why the $\delta^{(n)}(\nu)$-expansion for $f_V(\nu)$ starts with the $\delta'(\nu)$ term :

$$f_V(\nu) = \delta'(\nu) - L_V\,\delta''(\nu) + \dots , \tag{9}$$

with the parameter L_V determined by the magnitude of the condensates of dimension 8.

For the gluonic nonlocal condensate, in the Fock-Schwinger gauge, one has

$$\langle A_\mu^a(z)A_\nu^b(y)\rangle = \delta^{ab}\,(y_\mu z_\nu - g_{\mu\nu}(zy))\frac{\langle GG\rangle}{384}M_G((z-y)^2, z^2, y^2) + \dots , \tag{10}$$

where the M_G-function depends not only on the interval $(z-y)^2$, but also on z^2 and y^2. However, since the coefficients in front of z^2 and y^2 in the expansion

$$M_G = 1 - \frac{\langle GD^2G\rangle - \frac{2}{3}\langle j^2\rangle}{18\langle GG\rangle}\{(y-z)^2 + \frac{y^2+z^2}{8}\} + \dots \tag{11}$$

are rather small, one can start with the approximation

$$M_G(z^2, y^2, (z-y)^2) \simeq \int_0^\infty e^{\nu(z-y)^2/4} f_G(\nu)\,d\nu \tag{12}$$

introducing the distribution function $f_G(\nu)$.

There are three simplest trilocal quark-gluon condensates

$$\begin{aligned} M_{\mu\nu}(y,z) &\equiv \langle\bar{q}(0)\gamma_\nu A_\mu(y)q(z)\rangle = (z_\mu y_\nu - g_{\mu\nu}(zy))M_1 + (y_\mu y_\nu - g_{\mu\nu}y^2)M_2 + \dots && (13)\\ \tilde{M}_{\mu\nu}(y,z) &\equiv \langle\bar{q}(0)\gamma_\nu\gamma_5 A_\mu(y)q(z)\rangle = \epsilon_{\mu\nu\rho\sigma}y^\rho z^\sigma M_3 + \dots . && (14)\end{aligned}$$

The functions M_{1-3} can be parameterized by the triple integral representation:

$$M_i(z^2, y^2, (z-y)^2) = A_i\int_0^\infty e^{(\nu_1 z^2/4 + \nu_2 y^2/4 + \nu_3(z-y)^2/4)} f_i(\nu_1,\nu_2,\nu_3)\,d\nu_1 d\nu_2 d\nu_3. \tag{15}$$

where $A_i = \{-\frac{3}{2}A, 2A, \frac{3}{2}A\}$. The limiting case of the standard local condensates (corresponding to $\lambda_q^2 \to 0$) is obtained by the substitution $f_i(\nu_1,\nu_2,\nu_3) \to \delta(\nu_1)\delta(\nu_2)\delta(\nu_3)$.

Incorporating the nonlocal condensates as described above, one arrives at a modified diagram technique, with some lines and vertices being the ordinary perturbative ones, and some corresponding to the nonlocal condensates. Increasing the number of loops, one should consider the condensates containing more and more fields. We restrict our analysis here by the two-loop level. Then, in addition to those already listed, one encounters the four-quark condensate. To simplify the calculation, we apply the vacuum dominance hypothesis and factorize it into a product of two bilocal ones.

4.Sum rule. Using the representations (4)-(7), and calculating the coefficient functions we obtain a modified QCD sum rule, with the δ-functions of eq.(3) substituted by the functionals $\delta\Phi_i(x)$ of 6 vacuum distribution functions:

$$f_\pi^2\varphi_\pi(x) = \frac{M^2}{4\pi^2}(1-e^{-s_0/M^2})\Phi^{pert}(x) + \{4\bar{x}f_V(xM^2) + \sum_{i=1}^{4}\delta\Phi_i(x) + \delta\Phi_G(x)\} + (x \to \bar{x}) \quad (16)$$

where $\bar{x} = 1-x$, M^2 is the Borel parameter and

$$\Phi^{pert}(x) = 6x\bar{x}\{1 + C_F\frac{\alpha_s}{4\pi}[5 - \frac{\pi^2}{3} + \log^2(\frac{\bar{x}}{x})]\}$$

is the "perturbative" contribution (free quark loop plus $O(\alpha_s)$ radiative corrections).

The simplest contribution, proportional to the f_V-function taken at $\nu = xM^2$, is displayed explicitly in eq.(16). Other contributions have a more involved form (see ref.[17]).

The most intriguing conclusion to be drawn from eq.(16) is that $\varphi_\pi(x)$, *the longitudinal momentum distribution of quarks inside the pion, is directly related to* $f(\nu)$, *the virtuality distribution of quarks and gluons in the vacuum.* Therefore, it is very important to know the form of the latter to estimate the moments $\langle\xi^N\rangle_\pi$ for $N > 0$.

5.Modelling $f_i(\nu)$. To obtain the original SR (3), one should take the first term of the $\delta^{(n)}$–expansion for the $f(\nu)$'s. It should be understood that this approximation is really the simplest model for the distribution functions $f(\nu)$. However, such a model (used, as a matter of fact, by CZ [13]) is evidently too crude if the L_i-parameters characterizing the width of $f_i(\nu)$ are comparable in magnitude with the relevant hadronic scale. In this situation, instead of the standard expansion over the local condensates we propose to use an expansion in which the (relatively) large average virtuality of the vacuum fields is taken into account just in the first term. For the functions $M(z^2)$ having a finite widths of an order of μ^2, it is much more preferable to use the expansion of $f(\nu)$ over $\delta^{(n)}(\nu-\mu^2)$. The first term of this expansion

$$M(z^2) = M(0)\{\, e^{z^2\mu^2/4} + \ldots\} \quad (17)$$

takes into account the main effect caused by the finite width of the function $M(z^2)$, while subsequent terms describe effects due to the deviation of its form from the Gaussian one.

To construct the Gaussian ansätze one should know the second term of the z^2-expansion of the relevant nonlocal condensates, *e.g.*, incorporating eq.(4) we take $f_S(\nu) = \delta(\nu - \lambda_q^2/2)$.

For M_μ the situation is more complicated: L_V is determined by 5 different LC, the values of which are poorly known. The simplest model is to assume that all the nonlocal distributions have the same width. So, we take $f_V^{(1)}(\nu) = \delta'(\nu - \lambda_q^2/2)$. Of course, it is more reasonable to expect that the shift parameters L_i, though all of the same order of magnitude, are still numerically different. Another model for L_V is to extract the part proportional to $\langle\bar{q}D^2q\rangle\langle\bar{q}q\rangle$ from all the relevant LC of dimension 8 and neglect the remaining contributions. This gives the value $L_V = \frac{7}{20}\lambda_q^2$, rather close to the naive estimate.

In a similar way we construct the model for the trilocal functions:

$$f_i(\nu_1,\nu_2,\nu_3) = A_i\delta(\nu_1 - L_i^{(1)})\delta(\nu_2 - L_i^{(2)})\delta(\nu_3 - L_i^{(3)}). \tag{18}$$

One can try to determine $L_i^{(j)}$'s from the expansion of the relevant NLC by retaining only the $\langle\bar{q}D^2q\rangle\langle\bar{q}q\rangle$ part of the coefficients in front of z^2, y^2 or $(z-y)^2$, respectively. This gives $L_1^j = \{\frac{53}{288}, -\frac{1}{144}, \frac{2}{9}\}\lambda_q^2$ for the f_1–function, $L_2^j = \{\frac{1}{192}, \frac{517}{960}, \frac{31}{192}\}\lambda_q^2$ for the f_2–function and $L_3^j = \{-\frac{1}{32}, \frac{19}{72}, \frac{1}{6}\}\lambda_q^2$ for the f_3–function [17]. According to these estimates, the trilocal condensates in some directions decrease much slower than in the others, and sometimes even increase when the distance between the quarks increases, which is completely unrealistic. Hence, it is not safe to neglect other LC estimating the width parameters and, in the absence of a reliable model of the QCD vacuum, we simply assume that the trilocals decrease at the same rate in all directions and take $L_i^{(j)} = \lambda_q^2/2$.

To model the nonlocality effects for the gluonic contribution, we assume, by analogy with the quark case, that the $\delta(x)$ terms of the $O(\langle GG\rangle)$ contribution (eq.(3)) should be substituted by $\delta(x - L_G/M^2)$ in eq.(16), with $L_G = \frac{2}{9}\lambda_q^2$, as suggested by eq.(11).

6.Numerical estimates. Within the simplified version of our Gaussian ("delta-function") model for the nonlocal condensates, the pion wave function sum rule has the following form:

$$\begin{aligned}
f_\pi^2\varphi_\pi(x) = {} & \frac{M^2}{4\pi^2}(1 - e^{-s_0/M^2})\Phi^{pert}(x) + \frac{1}{24\pi}\alpha_s\langle GG\rangle\delta(x - \frac{2}{9}\Delta) + \\
& + \frac{8}{81M^4}\pi\alpha_s\langle\bar{q}q\rangle^2\Bigg\{\bar{x}\delta'(x-\Delta) + 18\frac{\theta(x<\Delta)}{\Delta^2(1-\Delta)}\bar{x}(x + (\Delta - x)\log(\bar{x})) + \\
& + \frac{3}{1-\Delta}\left[\frac{\delta(x-\Delta) - \delta(x-2\Delta)}{\Delta} - (1-\Delta)\delta(x-\Delta) + \frac{2}{3}(1-2\Delta)\frac{(2+\Delta)}{\Delta}\delta(x-2\Delta)\right] - \\
& - 2\bar{x}\frac{\theta(\Delta<x<2\Delta)}{\Delta}\left[\frac{3x}{1-\Delta} + \frac{2}{\Delta}\left((3 - \frac{\Delta+2\bar{x}}{1-\Delta}\right)\right]\Bigg\} + \\
& + (x \to \bar{x})
\end{aligned} \tag{19}$$

where $\Delta = \lambda_q^2/2M^2$.

Main observation is that in place of the $\delta(x)$-type contributions we have now either the δ-functions with the shifted arguments or the functions that are smooth at $x = 0$. In both cases, the moments of such terms decrease as N increases. Hence, for sufficiently large values of λ_q, there is no dramatic increase in the ratios of the condensate contributions to the perturbative term. Taking $\lambda_q^2 = 0.4\,GeV^2$ [15], we obtain for the lowest moments

$$\langle\xi^2\rangle = 0.25 \qquad \langle\xi^4\rangle = 0.12 \qquad \langle\xi^6\rangle = 0.07. \tag{20}$$

These values do not differ strongly from those corresponding to the asymptotic wave function. Therefore, it is not surprising that the model WF

$$\varphi_\pi^{mod,1}(x) = \frac{8}{\pi} f_\pi \sqrt{x(1-x)}, \qquad \varphi_\pi^{mod,2}(x) = 6 f_\pi x(1-x)\left(1 + \frac{8}{9}(1 - 5x(1-x))\right) \quad (21)$$

reproducing these values (20), are also close to the asymptotic wave function. The second model corresponds to the expansion over the Gegenbauer polynomials $C_n^{3/2}(\xi)$ (the eigenfunctions of the evolution equation [4, 6]).

Thus, the moments of the pion WF are rather sensitive to the functional form of the nonlocal condensates. The faster the NLC decrease with the distance, the faster is the decrease with N of the relevant contribution into the $\langle \xi^N \rangle$ sum rule. Of course, in the $\lambda_q \to 0$ limit, eq.(19) reduces to the original CZ sum rule (3),(1), and one obtains large CZ values for the moments. With $\lambda_q^2 = 0.4\, GeV^2$, the condensate terms still decrease more slowly with N than the perturbative contribution, and the $\langle \xi^N \rangle$-values (20) are still larger than $\langle \xi^N \rangle^{as}$. To get the asymptotic value for $\langle \xi^2 \rangle$, one should take $\lambda_q^2 = 1.2\, GeV^2$. Surprizingly enough, it is this huge value of λ_q^2 that is favoured by a calculation within a rather realistic QCD vacuum model developed by Shuryak [16]. The recent lattice result $\langle \xi^2 \rangle = 0.11$ [19], is still rather far from these values, but the disagreement might be essentially reduced by a renormalization factor (of order of 1.5) not included into the quoted lattice value.

Our results depend on the models we accepted for the nonlocal condensates. However, the sum rule is dominated by a single contribution (the second term in the braces in eq.(19)) which is due to the four-quark condensate $\langle \bar{q}(0)q(x)\bar{q}(y)q(z) \rangle$, factorized *via* the vacuum dominance hypothesis to the product of the simplest $\langle \bar{q}(0)q(z) \rangle$-type condensates. This factorization amounts to neglecting the dependence on the distance between the two $\bar{q}q$ pairs. If one takes this dependence into account, then the dominant term of eq.(19) will produce the contributions that will faster decrease with N, and the resulting $\langle \xi^N \rangle$ will be even farther from the CZ values.

7. Conclusions. Our basic idea in the present paper is that the nonperturbative information about the QCD vacuum structure should be accumulated in the functions describing the momentum distribution of the vacuum quark and gluonic fields. For the vacuum, these functions play the role analogous to that of the parton distributions in the case of the hadrons. Ideally, the vacuum distribution functions should be calculated from the theory of the QCD vacuum. In the absence of such a theory, one can incorporate the fact that the same vacuum distribution functions appear in different NLC-modified QCD sum rules for hadronic wave functions, parton distribution functions, hadronic form factors *etc.* This opens a possibility of finding the vacuum distribution functions (universal for all the hadrons) from the experimentally known hadronic functions.

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