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SINGLET AXIAL CONSTANT FROM QCD SUM RULES

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Белицкий А.В., Теряев О.В. E2-95-374
Синглетная аксиальная константа из правил сумм КХД

Анализируется синглетный аксиальный формфактор протона при малом переданном импульсе в методе правил сумм КХД с использованием нуклонного тока, в который явно введены глюонные степени свободы. Результатом работы является количественное предсказание синглетной аксиальной константы. Показано, что наиболее важную роль в анализе играют билокальные поправки.

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Belitsky A.V., Teryaev O.V. E2-95-374
Singlet Axial Constant from QCD Sum Rules

We analyze the singlet axial form factor of the proton for small momentum transferred in the framework of QCD sum rules using the interpolating nucleon current which explicitly accounts for the gluonic degrees of freedom. As the result we come to the quantitative prediction of the singlet axial constant. It is shown that the bilocal power corrections play the most important role in the analysis.

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1. In the last several years there has been an increasing interest in the deep inelastic structure function $g_1^p(x)$. It was provoked by the EMC result on the scattering of the longitudinally polarized muon beam on a longitudinally polarized hadron target. The unexpectedly small asymmetry found by EMC has led to the so called "spin crisis in the parton model" and has raised a number of questions about the understanding of the dynamics of the proton spin on the parton level, namely, how the nucleon spin is build up from the spins of its constituents. An enormous flood of theoretical investigations was generated in order to resolve the current "spin problem" [1]. The EMC measurement of the first moment of the polarized structure function Γ_1^p can via the Ellis-Jaffe sum rule [2]

$$\begin{aligned}\Gamma_1^p(Q^2) &= \int_0^1 dx g_1^p(x,Q^2)\\ &= \frac{1}{12}\left\{\left(G_A^{(3)}(0)+\frac{1}{\sqrt{3}}G_A^{(8)}(0)\right)\left(1-\left(\frac{\alpha_s}{\pi}\right)-3.5833\left(\frac{\alpha_s}{\pi}\right)^2-20.2153\left(\frac{\alpha_s}{\pi}\right)^3\right)\right.\\ &\left.+\frac{4}{3}G_A^{(0)}(0,Q^2)\left(1-\frac{1}{3}\left(\frac{\alpha_s}{\pi}\right)-1.0959\left(\frac{\alpha_s}{\pi}\right)^2\right)\right\} \qquad (1)\end{aligned}$$

be interpreted as a measurement of the singlet axial constant $G_A^{(0)}(0)$ of the proton. The last one turns out to be unexpectedly small in contradiction with the naïve parton model where it is fairly close to unity. The EMC reported for $G_A^{(0)}(0)$ the result which is compatible with zero. Two new experiments are under way to check their measurement of $g_1^p(x)$ and to measure the analogous neutron function $g_1^n(x)$. The recent analysis [3] of all proton data at $Q^2 = 10GeV^2$ gives $G_A^{(0)}(0) = 0.27\pm0.08\pm0.10$, still far from unity. So the problem reduces to the evaluation of $G_A^{(0)}(0)$ because the other two axial constants can be extracted reliably from the data on neutron and hyperon β-decays. In this paper we calculate it in the framework of QCD sum rule approach which till now seems to be the most powerful method for extraction of the information about the low energy properties of hadrons and the closest one to the first principles of the theory.

In the eq.(1) the functions $G_A^{(i)}(Q^2)$ are form factors at zero momentum transferred in the proton matrix elements of axial currents

$$\langle N(p_2,\lambda_2)|j_{5\mu}^{(i)}(0)|N(p_1,\lambda_1)\rangle = \bar{u}_N^{(\lambda_2)}(p_2)\left(G_A^{(i)}(Q^2)\gamma_\mu\gamma_5 - G_P^{(i)}(Q^2)q_\mu\gamma_5\right)u_N^{(\lambda_1)}(p_1), \qquad (2)$$

where i is a $SU(3)_f$ index, $q = p_2 - p_1$ and $Q^2 = -q^2$. There is an important difference between the behaviour of induced pseudoscalar form factors at small momentum q. Here the singlet pseudoscalar form factor does not acquire a Goldstone pole at $Q^2 = 0$, even in the chiral limit, contrary to the matrix elements of the octet currents. It is known that this limit, in which the masses of the three light quark flavours are neglected, is not far away from the real world of hadrons. In this limit there exist eight massless pseudoscalar mesons, serving as the Goldstone bosons. However the ninth pseudoscalar, the η'-meson, remains massive. This property will be used in the following in order to extract the value of $G_A^{(0)}(0)$ from the sum rules.

It was established [1] that the first moment Γ_1^p does not measure the contribution of the quark spins to the proton one. This happens because of anomalous non-conservation of the singlet axial current. For this reason we display the rôle of this profound feature of the theory from the very beginning exploiting the equation for the anomalous divergence[1]:

$$\partial_\mu j_{5\mu}^{(0)} = 2i\sum_q m_q \bar{q}\gamma_5 q - \left(\frac{N_f\alpha_s}{4\pi}\right) G^a_{\mu\nu}\widetilde{G}^a_{\mu\nu}, \quad (3)$$

where N_f is a number of flavours (later $N_f = 3$). Taking the divergence of eq.(2) for singlet axial current and making use of the last expression we come to the relation which directly connects in the chiral limit the nonforward matrix elements of gluon operator to the effective form factor $2m_N G^{(0)}_{eff}(Q^2) = 2m_N G^{(0)}_A(Q^2) + Q^2 G^{(0)}_P(Q^2)$ that is equal to the $2m_N G^{(0)}_A(0)$ at $Q^2 = 0$.

2. As a rule, all calculations of the nucleon characteristics use the particular three-quark current introduced by Ioffe [5]. When one makes an attempt to evaluate the matrix elements of quark-gluon or gluon operators one faces the evident calculational difficulties, moreover the final sum rules are aggravated by extra UV logarithms due to mixing of operators and therefore the calculations are affected by noncontrollable uncertainties [6].

In field theory, the usual statement that the nucleon consists mainly of the three quarks means that the *3 quarks → 3 quarks* Green function (three quarks are in a state with nucleon quantum numbers) has the pole at the mass of the nucleon, with total angular momentum $J = \frac{1}{2}$, with a significant residue. The fact that the nucleon is not just a three quark state means that the nucleon pole also occurs, albeit with smaller residue, in a Green functions such as *3 quarks+g → 3 quarks+g*. For this reason one is forced to introduce a more complicated interpolating proton field which explicitly contains the gluonic degrees of freedom:

$$\eta_G(x) = \epsilon^{ijk}\left(u^i(x)C\gamma_\mu u^j(x)\right)\gamma_5\gamma_\mu\sigma_{\alpha\beta}\left(gG^a_{\alpha\beta}(x)t^a d(x)\right)^k. \quad (4)$$

The latter was investigated in ref. [7] and checked in the calculation of proton gluonic form factor normalized to the fraction of nucleon momentum carried by gluons. Recently, making use of this current the twist-3 and twist-4 corrections to Bjorken and Ellis-Jaffe sum rule were found [8]. The advantages of this current are straightforward: the calculations are simplified drastically, sum rules are free from the additional divergences that are not removed by the Borel transformation.

In the usual approach the nucleon matrix element of local operator can be extracted from the three-point correlation function which can be appropriately decomposed into several tensor structures. In our case it looks like:

$$\begin{aligned} W(p_1,p_2,q) &= i^2\int d^4x d^4y e^{ip_1x - ip_2y}\langle 0|\eta_G(x)\left(\frac{N_f\alpha_s}{4\pi}\right)G^a_{\mu\nu}(0)\widetilde{G}^a_{\mu\nu}(0)\,\bar{\eta}_G(y)|0\rangle \\ &= \sigma_{\mu\nu}\gamma_5 p_{1\mu}p_{2\nu}W^{(1)}(p_1^2,p_2^2,q^2) + \not{q}\gamma_5 W^{(2)}(p_1^2,p_2^2,q^2) + q^2\gamma_5 W^{(3)}(p_1^2,p_2^2,q^2). \end{aligned} \quad (5)$$

In practical calculation it is advantageous to consider the $W^{(1)}(p_1^2,p_2^2,q^2)$ because of its lower dimensionality. Another reason in favour of this choice is that it does not lead to the fictitious kinematical singularities in q^2 as the last term in the eq.(5) does. For this invariant amplitude we can write the double dispersion representation:

$$W^{(1)}(p_1^2,p_2^2,q^2) = \frac{1}{\pi^2}\int_0^\infty\int_0^\infty ds_1 ds_2\frac{\rho(s_1,s_2,Q^2)}{(s_1-p_1^2)(s_2-p_2^2)} + \ldots, \quad (6)$$

[1]Throughout the paper, we adopt the conventions in Itzykson and Zuber [4].

where the ellipses stand for the polynomials in p_1^2 and p_2^2 which die out after the double Borel transformation has been applied. For the physical spectral density we have accepted the conventional "resonance plus continuum" model:

$$\begin{aligned} \rho(s_1,s_2,Q^2) &= \pi^2 m_N^4\lambda_G^2 2m_N G^{(0)}_{eff}(Q^2)\delta(s_1-m_N^2)\delta(s_2-m_N^2) \\ &+ \rho_{cont}(s_1,s_2,Q^2)(1-\theta(\sigma_0-s_1)\theta(\sigma_0-s_2)). \end{aligned} \quad (7)$$

The function in front of double-pole term is a combination of form factors we are interested in up to certain overlap λ_G between the state created from the vacuum by η_G and the nucleon state:

$$\langle 0|\eta_G(0)|N(p,\lambda)\rangle = m_N^{2} \lambda_G u_N^{(\lambda)}(p). \quad (8)$$

So, our aim is the evaluation of the correlation function (5) in QCD. In the case when all the momenta $(-p_1^2) \sim (-p_2^2) \sim Q^2$ are sufficiently large (of order of $1GeV^2$) the leading contribution comes from the domain where all distances are small. Thus the standard machinery of short distance expansion are applicable, allowing to express the final result in terms of quark and gluon condensates. The problem modifies drastically if the squared momentum transferred become small ($Q^2 \ll (-p_i^2)$) because the relevant t-channel distances can be large. In this case the OPE has a twofold structure [9]. Terms of the first type arise from the SD(I)-region when all intervals $x^2 \sim y^2 \sim (x-y)^2$ are small. Another contribution comes from SD(II)-region (bilocal power correction) which originate from distances $x^2 \sim y^2 \gg (x-y)^2$. The necessity for the bilocal power corrections can be traced from the fact that the ordinary QCD Feynman diagrams contributing to the form factor at moderately large Q^2 in the limit of small Q^2 posses the logarithmical non-analiticities $(Q^2)^n lnQ^2$ which signals that the large distances come into play [10]. Therefore we have to subtract such *perturbative* behaviour from the corresponding graphs and add the "exact" correlators which account for the nonperturbative effects and thus possess the correct analytical properties as Q^2 goes to zero [11].

The large distance contribution is determined by the two-point correlator of operator of interest and some nonlocal string operator with definite twist (not dimension) which arises from the OPE of T-product of nucleon currents. The bilocal power corrections cannot be directly calculated in perturbation theory but we can write down the dispersion relation for them:

$$W(q^2,(xq),x^2) = \frac{1}{\pi}\int_0^\infty ds\frac{\rho(s,(xq),x^2)}{s-q^2}, \quad (9)$$

assuming the standard spectral density model with continuum starting at some threshold s_0 and finding in some way its parameters. We always do this constructing the auxiliary sum rules. There is no need in additional subtractions in eq.(9) because one always deals with the difference between the "exact" bilocal and its perturbative part, so due to coincidence of their UV behaviours the subtraction terms cancel in this difference.

Calculating the bilocal power corrections we account for the lowest twist-3 operators because we expect the higher twist contributions to be rather small. As can be easily traced from the eq.(10) such a contribution arises only (in the symmetrical kinematics $\tau_1 = \tau_2$ which we will accept in the following for the Borel parameters) for bilocals with the coefficient function containing the quark condensate. All other bilocals correspond to the next-to-leading twists and thus will be omitted below.

Collecting all contributions we come to the Borel sum rule with the Borel parameters τ_1 and τ_2:

$$\begin{aligned}
&\frac{m_N^4\lambda_G^2}{\tau_1\tau_2}2m_N G^{(0)}_{eff}(Q^2)e^{-\frac{m_N^2}{\tau_1}-\frac{m_N^2}{\tau_2}} \\
&= -\frac{1}{\pi^2}\int_0^\infty\int_0^\infty \frac{ds_1 ds_2}{\tau_1\tau_2}(1-\theta(\sigma_0-s_1)\theta(\sigma_0-s_2))\rho_{cont}(s_1,s_2,Q^2)e^{-\frac{s_1}{\tau_1}-\frac{s_2}{\tau_2}} \\
&+\frac{1}{\pi^2}\frac{N_f}{3}\left(\frac{\alpha_s}{\pi}\right)^2\langle\bar uu\rangle\frac{(\tau_1\tau_2)^2}{(\tau_1+\tau_2)^2}\left\{\frac{Q^2}{(\tau_1+\tau_2)}J_{12}+\left(I_4+I_6+\frac{1}{6}I_8\right)e^{\frac{Q^2}{(\tau_1+\tau_2)}}\right\} \\
&-\frac{1}{\pi^2}\frac{N_f}{6}\left(\frac{\alpha_s}{\pi}\right)^2 m_0^2\langle\bar uu\rangle\frac{(\tau_1\tau_2)}{(\tau_1+\tau_2)}\left\{\frac{Q^2}{(\tau_1+\tau_2)}J_{11}+\left(I_4+\frac{1}{2}I_6\right)e^{\frac{Q^2}{(\tau_1+\tau_2)}}\right\} \\
&+\frac{1}{\pi^2}\frac{N_f}{144}\left(\frac{\alpha_s}{\pi}\right)^2 m_0^2\langle\bar uu\rangle\left\{4\frac{(\tau_1\tau_2)^2}{(\tau_1+\tau_2)^3}\left[J_{12}+\left(I_2+I_4+\frac{1}{6}I_6\right)e^{\frac{Q^2}{(\tau_1+\tau_2)}}\right]\right. \\
&\qquad\left.+\frac{(\tau_1\tau_2)}{(\tau_1+\tau_2)}\left[\frac{Q^2}{(\tau_1+\tau_2)}J_{02}-\left(I_2+2I_4+\frac{1}{2}I_6\right)e^{\frac{Q^2}{(\tau_1+\tau_2)}}\right]\right\} \\
&-\frac{1}{\pi^2}\frac{7N_f}{32}\left(\frac{\alpha_s}{\pi}\right)^2 m_0^2\langle\bar uu\rangle\frac{(\tau_1\tau_2)^2}{(\tau_1+\tau_2)^3}\left\{J_{12}+\left(I_2+I_4+\frac{1}{6}I_6\right)e^{\frac{Q^2}{(\tau_1+\tau_2)}}\right\} \\
&-\frac{1}{\pi^2}\frac{N_f}{8}\left(\frac{\alpha_s}{\pi}\right)^2 m_0^2\langle\bar uu\rangle\left\{2\frac{(\tau_1\tau_2)^2}{(\tau_1+\tau_2)^3}\left[J_{12}+\left(I_2+I_4+\frac{1}{6}I_6\right)e^{\frac{Q^2}{(\tau_1+\tau_2)}}\right]\right. \\
&\qquad\left.-\frac{(\tau_1\tau_2)}{(\tau_1+\tau_2)}\left[\frac{Q^2}{(\tau_1+\tau_2)}J_{02}-\left(I_2+2I_4+\frac{1}{2}I_6\right)e^{\frac{Q^2}{(\tau_1+\tau_2)}}\right]\right\} \\
&-\frac{16}{3N_f}\langle\bar uu\rangle\frac{(\tau_1\tau_2)}{(\tau_1+\tau_2)^2}e^{\frac{Q^2}{(\tau_1+\tau_2)}}\frac{m_{\eta'}^2 f_{\eta'}}{m_{\eta'}^2+Q^2}\left(f^{(1)}_{\eta'}\varphi^{(1)}_{\eta'}\left(\frac{\tau_1}{\tau_1+\tau_2}\right)+f^{(2)}_{\eta'}\varphi^{(2)}_{\eta'}\left(\frac{\tau_1}{\tau_1+\tau_2}\right)\right),
\end{aligned}\tag{10}$$

where the continuum double spectral density is

$$\rho_{cont}(s_1,s_2,Q^2)=\sum_{i=1}^{5}\rho_{(i)}(s_1,s_2,Q^2). \tag{11}$$

and each term in a sum is found from the corresponding diagram in fig.1:

$$\rho_{(1)}(s_1,s_2,Q^2)=\frac{N_f}{72}\left(\frac{\alpha_s}{\pi}\right)^2\langle\bar uu\rangle Q^4\left(1-\frac{\sigma}{R^{\frac12}}\right)^2\left(2+\frac{\sigma}{R^{\frac12}}\right), \tag{12}$$

$$\rho_{(2)}(s_1,s_2,Q^2)=-\frac{N_f}{6}\left(\frac{\alpha_s}{\pi}\right)^2 m_0^2\langle\bar uu\rangle Q^4\frac{s_1s_2}{R^{\frac32}}, \tag{13}$$

$$\begin{aligned}
&\rho_{(3)}(s_1,s_2,Q^2) \\
&=\frac{N_f}{144}\left(\frac{\alpha_s}{\pi}\right)^2 m_0^2\langle\bar uu\rangle Q^2\left\{\frac{1}{6}\left(1-\frac{\sigma}{R^{\frac12}}\right)^2\left(2+\frac{\sigma}{R^{\frac12}}\right)-\frac{1}{2}\left(1-\frac{\sigma}{R^{\frac12}}\right)-Q^2\frac{s_1s_2}{R^{\frac32}}\right\},
\end{aligned}\tag{14}$$

$$\rho_{(4)}(s_1,s_2,Q^2)=-\frac{7N_f}{768}\left(\frac{\alpha_s}{\pi}\right)^2 m_0^2\langle\bar uu\rangle Q^2\left(1-\frac{\sigma}{R^{\frac12}}\right)^2\left(2+\frac{\sigma}{R^{\frac12}}\right), \tag{15}$$

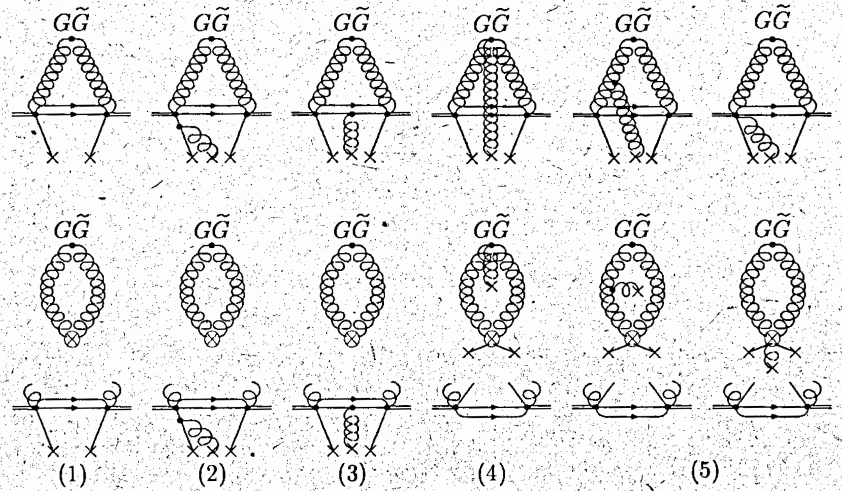

Fig.1.Contributions to the effective axial form factor in QCD sum rule approach.

$$\begin{aligned}
&\rho_{(5)}(s_1,s_2,Q^2) \\
&=-\frac{N_f}{8}\left(\frac{\alpha_s}{\pi}\right)^2 m_0^2\langle\bar uu\rangle Q^2\left\{\frac{1}{12}\left(1-\frac{\sigma}{R^{\frac12}}\right)^2\left(2+\frac{\sigma}{R^{\frac12}}\right)+\frac{1}{2}\left(1-\frac{\sigma}{R^{\frac12}}\right)+Q^2\frac{s_1s_2}{R^{\frac32}}\right\},
\end{aligned}\tag{16}$$

and

$$\sigma=s_1+s_2+Q^2,\qquad R(s_1,s_2,Q^2)=\sigma^2-4s_1s_2. \tag{17}$$

The functions J_{nm} are originated from the diagrams in the first row in fig.1 and are given by the following expression:

$$J_{nm}(\tau_i,Q^2)=\int_0^1 dx\bar x^{n-1}x^m exp\left\{-\frac{x}{\bar x}\frac{Q^2}{(\tau_1+\tau_2)}\right\} \tag{18}$$

The functions I_n represent the difference between the perturbative part of the bilocal power corrections (the second row in fig.1) and the continuum contribution into the "exact" bilocals:

$$I_n(\tau_i,Q^2,s_0)=(\tau_1+\tau_2)^{-n/2}\left[(Q^2)^{n/2}ln\left(\frac{s_0+Q^2}{Q^2}\right)+\sum_{k=1}^{n/2}\frac{(-1)^k}{k}s_0^k(Q^2)^{n/2-k}\right], \tag{19}$$

As can be easily seen the expressions in the brackets in eq.(10) are free from logarithmic non-analiticities, they are replaced by the combination s_0+Q^2 which are "safe" in the limit $Q^2\to 0$.

We state that contrary to the refs. [8] where the sum rules with the same interpolating nucleon field were dominated by the contribution from the highest dimension operators our sum rule does not affected by them: the coefficient functions that are determined to

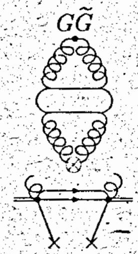

Fig.2.Leading twist bilocal power correction.

the leading accuracy by tree and one-loop diagrams vanish identically. Therefore we do not meet the problem of breakdown of OPE for the correlator in question. The last line in eq.(10) represents the leading bilocal correction, see the graph in fig.2. The functions $\varphi^{(i)}$ are the wave functions describing the light-cone momentum distribution of gluon inside meson. In the following we will take for them the asymptotical form which seems not far away from the reality:

$$\varphi^{(1)}(\alpha) = 30\alpha^2\bar{\alpha}^2,$$
$$\varphi^{(2)}(\alpha) = 420(\alpha-\bar{\alpha})\alpha^2\bar{\alpha}^2 \tag{20}$$

The unknown overlaps $f_{\eta'}$, $f_{\eta'}^{(1)}$ and $f_{\eta'}^{(2)}$ of some gluon operators with the η'-meson state are defined as:

$$\langle 0|\frac{N_f\alpha_s}{4\pi}G^a_{\mu\nu}(0)\widetilde{G}^a_{\mu\nu}(0)|\eta'(q)\rangle = m^2_{\eta'}f_{\eta'}$$
$$\langle\eta'(q)|\frac{N_f\alpha_s}{4\pi}\Big(G^a_{\mu\lambda}(x)\widetilde{G}^a_{\mu\kappa}(0)-G^a_{\mu\kappa}(x)\widetilde{G}^a_{\mu\lambda}(0)\Big)|0\rangle$$
$$= (x_\kappa q_\lambda - x_\lambda q_\kappa)[i f^{(1)}_{\eta'}\phi^{(1)}_{\eta'}(xq)+f^{(2)}_{\eta'}(xq)\phi^{(2)}_{\eta'}(xq)]. \tag{21}$$

and may be found from the additional sum rules. In the last line the wave functions $\phi^{(i)}_{\eta'}$ can be related in the standard way to the usual momentum fraction ones:

$$\phi^{(i)}(xq) = \int_0^1 d\alpha e^{i\alpha(xq)}\varphi^{(i)}(\alpha). \tag{22}$$

The $f_{\eta'}$ constant was calculated in ref. [12] in connection with pseudoscalar gluonium, while the other two will be extracted below.

3. We estimate the bilocal power-correction saturating the two-point correlator:

$$\Pi_{\lambda\kappa}(x;q)$$
$$= i\int d^4y e^{iqy}\langle 0|T\left\{\frac{N_f\alpha_s}{4\pi}G^a_{\mu\nu}(y)\widetilde{G}^a_{\mu\nu}(y)\ \frac{N_f\alpha_s}{4\pi}\Big(G^a_{\mu\lambda}(x)\widetilde{G}^a_{\mu\kappa}(0)-G^a_{\mu\kappa}(x)\widetilde{G}^a_{\mu\lambda}(0)\Big)\right\}|0\rangle \tag{23}$$

by the contribution of low lying resonances with corresponding quantum numbers, in our case by η'-meson. It is likely to be the only prominent *singlet* pseudoscalar both in quark and gluon channels. It turns out that due to the antisymmetrical tensor structure

involved the contribution of ordinary local power corrections with gluon condensates vanish identically in the theoretical part of the sum rules. For this reason we account for nonperturbative effects introducing the concept of nonlocal gluon condensate [13] which corresponds to infinite series of local ones. It can be appropriately decomposed into two tensor structures times corresponding form factors [14]:

$$\langle 0|G^a_{\mu\rho}(x)\widetilde{E}^{ab}(x,0)G^b_{\nu\sigma}(0)|0\rangle = \frac{\langle G^2\rangle}{12}\Big\{(g_{\mu\nu}g_{\rho\sigma}-g_{\mu\sigma}g_{\nu\rho})\left(D_{NA}(x^2)+D_A(x^2)\right)$$
$$+(g_{\mu\nu}x_\rho x_\sigma + g_{\rho\sigma}x_\mu x_\nu - g_{\mu\sigma}x_\nu x_\rho - g_{\nu\rho}x_\mu x_\sigma)\frac{dD_A(x^2)}{dx^2}\Big\}. \tag{24}$$

This form explicitly separates out the term proportional to $D_{NA}(x^2)$ which violate the abelian Bianchi identity, while the second term satisfies it. It was shown that linear confinement occurs when $D_{NA}(x^2)$ is present in (24) while the second term does not contribute to the string tension [14].

In the calculation of $f^{(1)}_{\eta'}$ and $f^{(2)}_{\eta'}$ constants only abelian form factor contribute. We present it in the form of α-representation [15]:

$$D_A(x^2) = \int_0^\infty d\alpha f^A_G(\alpha,\lambda^2_A)e^{\alpha\frac{x^2}{4}}. \tag{25}$$

and use a δ-shaped ansatz for the distribution function $f^A_G(\alpha,\lambda^2_A)$:

$$f^A_G(\alpha,\lambda^2_A) = \delta(\alpha-\lambda^2_A), \tag{26}$$

where $1/\lambda_A$ is an abelian correlation length of the vacuum fluctuations, it can be expressed in terms of vacuum condensates $\lambda^2_A = \frac{8}{9}g^2\langle\bar{u}u\rangle^2/\langle G^2\rangle \approx 0.03GeV^2$ at $1GeV^2$ [16].

Proceeding in the standard way we obtain the following sum rule:

$$m^2_{\eta'}f_{\eta'}f^{(1)}_{\eta'}e^{-\frac{m^2_{\eta'}}{M^2}} = M^2\frac{N_f\alpha_s}{4\pi}\left(\frac{1}{3\pi^2}\frac{N_f\alpha_s}{4\pi}M^4E_2\left(\frac{s_0}{M^2}\right)+\frac{N_f}{3}\langle\frac{\alpha_s}{\pi}G^2\rangle\frac{\lambda^2_A}{M^2}\left(1-\frac{\lambda^2_A}{M^2}\right)\right), \tag{27}$$

where

$$E_2(x) = 1-(1+x+\frac{x^2}{2})e^{-x}, \tag{28}$$

while the ratio $|f^{(2)}_{\eta'}/f^{(1)}_{\eta'}| \sim 1/140$ turns out very small and we neglect the second constant in the following analysis.

We stress that due to the fact that one of the gluon currents has nonzero Lorentz spin leads to the absence of direct instantons to the polarization operator of interest [17]. In the correlator considered the power correction are small and we see no reason why the local duality has to fail. Then taking the limit $M^2 \to \infty$ we obtain:

$$m^2_{\eta'}f_{\eta'}f^{(1)}_{\eta'} = \frac{N_f\alpha_s}{4\pi}\left(\frac{1}{3\pi^2}\frac{N_f\alpha_s}{4\pi}\frac{s_0^3}{6}+\frac{N_f}{3}\langle\frac{\alpha_s}{\pi}G^2\rangle\lambda^2_A\right). \tag{29}$$

The value of continuum threshold is taken from the analysis performed in [18]: $s_0 = 2.5GeV^2$.

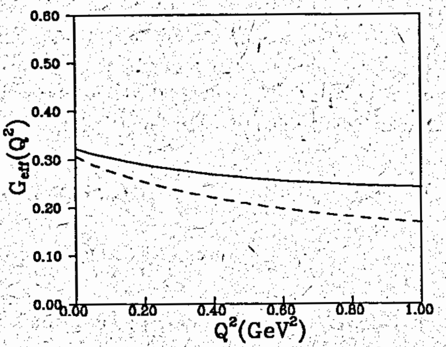

Fig.3. Effective axial form factor normalized to the singlet axial constant: QCD sum rule prediction for $\tau = 1.2GeV^2$ shown by solid curve and for $\tau = 1.7GeV^2$ by dashed curve.

We use the standard ITEP values of condensates rescaled to the normalization point $\mu^2 \sim m_N^2 \sim 1GeV^2$ with the appropriate anomalous dimensions: $\langle \bar{u}u \rangle = (-0.257GeV)^3$, $m_0^2 = \langle \bar{u}g(\sigma G)u \rangle / \langle \bar{u}u \rangle = 0.65GeV^2$ and continuum threshold $\sigma_0 = (1.5GeV)^2$.

The resulting behaviour of the form factor is shown in fig.3 corresponding to two values of Borel parameter $\tau = \tau_1 = \tau_2$: solid and dashed curves correspond to the values $\tau = 1.2GeV^2$ and $\tau = 1.7GeV^2$ respectively.

4. Because the axial constant is not a renormgroup invariant we have to evolve it from QCD sum rule scale $\mu^2 \sim 1GeV^2$ up to the one of EMC-SMC experiment which is $Q^2 = 10GeV^2$ exploiting the one-loop solution of RG equation:

$$G_A^{(0)}(0,Q^2) = G_A^{(0)}(0,\mu^2) exp\left\{\frac{\gamma_2}{4\pi\beta_0}[\alpha_s(Q^2) - \alpha_s(\mu^2)]\right\}, \tag{30}$$

where the anomalous dimension $\gamma_2 = 16N_f$ and as usual $\beta_0 = 11 - \frac{2}{3}N_f$. Thus we obtain

$$G_A^{(0)}(0, Q^2 = 10GeV^2) = 0.29, \tag{31}$$

with the accuracy not to be worse than 30%. Errors are due to the τ dependence of the sum rule, the uncertainty in the value of continuum threshold s_0 and the bilocal power corrections omitted. This value are in good agreement with the new world average value for the singlet constant.

Numerically it is the bilocal power correction that the most important, while all dependence on the particular baryon current are absorbed into coefficient function. With the fact on hand we come to the same conclusion as the authors of ref. [19] that the suppression of the flavour singlet component of the first moment of g_1^p observed by EMC-SMC collaboration is a target-independent feature of QCD and is not the property of the proton structure.

In ref.[18] an attempt was undertaken to evaluate $G_A^{(0)}(0)$ by QCD sum rules in a way similar to the calculation of the octet axial constant [20]. However, due to the presence of the anomaly in the induced vacuum condensates the problem differs significantly from the one for the $G_A^{(8)}(0)$. This feature was incorporated in the calculation but nevertheless the authors did not come to the reasonable quantitative prediction of the singlet axial constant. And it was conjectured that the OPE breaks down for the singlet axial current in the *axial-nucleon-nucleon* vertex. However, we did not observe any evidence of the divergence of the OPE in the correlator under investigation: the contribution of the highest dimension vacuum condansates unsuppressed by a number of loops is absent.

A possible line of development would be to estimate twist-three gluon contribution into the moments of the transverse spin structure function g_2 [21], twist-two polarized gluon distribution in the nucleon. However, the latter would require an elaboration of the procedure for seperation of the large and small distances in the effective four-point correlator.

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