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V.M. Khatsymovsky, I.B. Khriplovich
A.R. Zhitnitsky

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V.M.Khatsymovsky, I.B.Khriplovich, A.R.Zhitnitsky

Institute of Nuclear Physics, 630090, Novosibirsk, U.S.S.R.

Abstract

The arguments in favour of large value of the matrix element $\langle N|\bar{s}s|N\rangle$ are given which are based on some semiphenomenological data and low-energy theorems. As an application of presented considerations we study the contribution of strange quarks to the neutron electric dipole moment.

1. Introduction

The point of view is widely spread that the admixture of the pairs of strange quarks $\bar{s}s$ in nucleons is small. Deep inelastic scattering data indicate to this smallness. It should be noted however that this argument is valid for s--quarks in the vector channel only. As for the scalar channel, there is no serious reasons to believe that $\langle N|\bar{s}s|N\rangle$ is much more smaller than corresponding $\bar{u}u$ and $\bar{d}d$ matrix elements. In recent work [1] the arguments were given indicating that the matrix element $\langle N|\bar{s}s|N\rangle$ is not small indeed. This conclusion seems to us to be very interesting by itself.

In this work we present additional arguments based on low-energy theorems and some semiphenomenological data in favour of large nucleon expectation value of $\bar{s}s$. A similar approach being applied to calculation of baryon matrix elements of (pseudo-)scalar operators allows one to make important conclusions on the interaction of Higgs particle with nucleons and also on CP-odd effects in the system of nonstrange baryons.

In particular, prediction of Weinberg model of CP-violation for the nucleon electric dipole moment contradicts definitely the experimental data.

2. Quark scalar expectation

Let us start with calculation of quark scalar matrix elements over nucleon under the assumption of octet nature of SU(3) symmetry breaking. This calculation differs from that of ref. [1], section 2, by technical details only. However we present it here for completeness.

Averaging the results of various fits to the data on πN -scattering presented in ref. [2] leads to the following answer for the so-called σ -term

$$\frac{m_u+m_d}{2}\langle p|\bar{u}u+\bar{d}d|p\rangle = 63\pm 12\,MeV \quad (1)$$

(Here and below we omit kinematical structures like $\bar{p}p$ in the expressions for matrix elements).

Taking the values of quark masses to be m_u = 4 MeV, m_d = 7 MeV we get from (1)

$$\langle p|\bar{u}u+\bar{d}d|p\rangle = 11,5\pm 2 \quad (2)$$

Further, assuming octet type of SU(3) breaking responsible for mass splittings in the baryon octet we find

$$\langle p|\bar{u}u-\bar{d}d|p\rangle = \frac{m_\Xi - m_\Sigma}{m_s} = 0,9 \quad (3)$$

$$\langle p|\bar{u}u+\bar{d}d-2\bar{s}s|p\rangle = 3\,\frac{m_\Xi - m_\Lambda}{m_s} = 4,3 \quad (4)$$

Here m_Ξ, m_Σ, m_Λ are the masses of Ξ, Σ, and Λ hyperons respectively; for s-quark mass we take the value m_s = 140 MeV.

Using formulas (2)-(4) it is not difficult to obtain

$$\langle p|\bar{u}u|p\rangle = 6,2 \quad (5)$$

$$\langle p|\bar{d}d|p\rangle = 5,3 \quad (6)$$

$$\langle p|\bar{s}s|p\rangle = 3,6 \quad (7)$$

It should be mentioned that the accuracy of the relations (5)-(7) is not too high. The error in the value of the σ-term (see (1), (2)) already leads to the error of the order of 1 in each of the matrix elements found.

We would like to present at once the following very simple (and, to our point of view, quite convincing) argument in favour of the relations (5)-(7). It is clear that the average (7) is connected with the see quarks and is due to vacuum effects. It is natural to suppose that the vacuum contribution

to the averages (5), (6) is in the same ratio to the value (7) as the vacuum averages are:

$$\langle 0|\bar{u}u|0\rangle : \langle 0|\bar{s}s|0\rangle \simeq 1,2$$

Then the contribution of valence quarks is

$$\langle p|\bar{u}u|p\rangle_1 = 1,9 \quad (5a)$$

$$\langle p|\bar{d}d|p\rangle_1 = 1,0 \quad (6a)$$

These values are in remarkable agreement with the numbers 2 and 1, which could be expected from the naive picture of non-relativistic constituent quarks.

Thus together with the authors of ref. [1] we believe that the resolution of the known contradiction [2,3] between the experimental value of σ - term and assumptions of octet picture of SU(3) breaking and of smallness of the expectation value $\langle p|\bar{s}s|p\rangle$ consists in this matrix element being by no means small, in accordance with (7).

Now let us give some more arguments in favour of this assertion which supplement those given above and those based on Skyrme model and bag model given in ref. [1].

3. $0^\pm$ channels and their role in the nucleon physics

In this section we present arguments in favour of large value of $\langle N|\bar{s}s|N\rangle$ based on some low-energy theorems, QCD sum rules and phenomenological results. We proceed from the statement of [4] that (pseudo-)scalar channel is distinguished from others by it's strong coupling to the vacuum.

To study the problem by means of QCD sum rules [5,6,7] consider the correlator

$$T(p^2) = \int e^{ipx}\,dx\,dy\,\langle 0|T\{\eta(x)\bar{s}s(y)\bar{\eta}(0)\}|0\rangle \quad (8)$$

at $-p^2\to\infty$. Here η is an arbitrary nucleon current*. Note that we pick out the unit matrix kinematical structure in (8). Due to the absence of S -quark field in the nucleon current η any substantial contribution to $T(p^2)$ is connected only with nonperturbative correlators ("induced vacuum expectation values") [7] of the type

$$K = i\int dy \langle 0|T\{\bar{s}s(y)\bar{u}u(0)\}|0\rangle, \tag{9}$$

see fig. 1. As usually restricting to the lowest nucleon states in (8) we have

$$T(p^2) = \beta\frac{i}{\hat{p}-m}\alpha\frac{i}{\hat{p}-m}\bar{\beta} \tag{10}$$

where $\beta = \langle 0|\eta|N\rangle$, $\bar{\beta} = \langle N|\bar{\eta}|0\rangle$, $\alpha = \langle N|\bar{S}S|N\rangle$
On the other hand, the same β, $\bar{\beta}$ are defined by the polarization operator

$$\Pi(p^2) = \int e^{ipx} dx \langle 0|T\{\eta(x)\bar{\eta}(0)\}|0\rangle = \beta\frac{i}{\hat{p}-m}\bar{\beta} \tag{11}$$

The main contribution to the unit matrix structure (the same as in (8)) is given by the graph of fig. 2 analogous to that of fig. 1 up to the substitution of $i\langle\bar{q}q\rangle$ for K. Comparing (10) with (11) at $p^2 \ll -1\,GeV^2$ we get

$$\langle N|\bar{S}S|N\rangle \simeq m\frac{-1}{\langle\bar{q}q\rangle}K \tag{12}$$

Thus, calculation of $\langle N|\bar{S}S|N\rangle$ reduces to the evaluation of the vacuum correlator K. Fortunately, we get sufficiently rich information about the latter both from the low-energy theorems [4] and from the phenomenological analysis.

Digressing for a moment from the calculation of $\langle N|\bar{S}S|N\rangle$

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* The current may be chosen, e.g., in the standard form $\eta = \varepsilon^{abc}\gamma^{\mu}d^{a}(u^{b}C\gamma_{\mu}u^{c})$ [7]; note, however, that the results obtained below do not imply such the concretization.

note that this method of reduction of the nucleon matrix element problem to that of vacuum correlator is directly generalized to cover arbitrary scalar $\mathcal{O}_S$ or pseudoscalar $\mathcal{O}_P$ operators:

$$\langle B|\mathcal{O}_S|B\rangle \simeq \frac{m\bar{B}B}{-\langle\bar{q}q\rangle} i\int dy \langle 0|T\{\mathcal{O}_S(y)\bar{q}q(0)\}|0\rangle \tag{13a}$$

$$\langle B|\mathcal{O}_P|B\rangle \simeq \frac{m\bar{B}\gamma_5 B}{-\langle\bar{q}q\rangle} i\int dy \langle 0|T\{\mathcal{O}_P(y)\bar{q}\gamma_5 q(0)\}|0\rangle \tag{13b}$$

Also note that if there are the fields which present both in $\mathcal{O}_S$ (or $\mathcal{O}_P$) and in η , the contribution of small distances (due to, e.g., the diagrams 3) must be taken into account as well.

It will be shown in the framework of this approach that the value $\langle N|\bar{S}S|N\rangle$ is large and comparable to (7). How does it agree with the known smallness of $\langle N|\bar{S}\gamma_\mu S|N\rangle$? The answer is that the S -quarks in the nonvacuum channel do not turn into u, d -quarks: the Zweig rule works well. Phenomenologically it shows up in, e.g., the smallness of $\varphi - \omega$ -mixing [5] . Correspondingly correlators of the type $\int dy \langle 0|T\{\bar{S}\gamma_\mu S(y)\bar{u}\gamma_\nu u(0)\}|0\rangle$ are numerically small. Things are different in the scalar channel: the Zweig rule is badly broken and there is substantial admixture of S -quarks in the scalar mesons S^*, δ, ε and in (both) η and η'. The matter is that the correlators of the type (9) are of the same order of magnitude as $i\int dy \langle 0|T\{\bar{u}u(y)\bar{u}u(0)\}|0\rangle$ is. To summarize, the large value of (7) and violation of Zweig rule in the scalar channel are the manifestations of the same physics, namely that of scalar vacuum channel.

One more remark. Let us calculate the matrix element $\langle N|\bar{u}u|N\rangle$ by means of sum rules. In this case contribution from small distances $\sim m^{-1}$ (the diagrams like that shown in fig. 3) accounts for less than 30% of experimental value (2). Naively, fig. 3 corresponds to the valence quark contribution to the matrix element $\langle N|\bar{u}u|N\rangle$. From our point of view the rest is due to the vacuum contributions (see fig. 1) but not to the valence ones. The estimates of the correlator (9) given below confirm this conjecture.

The simplest way to evaluate (9) is as follows. Consider the low-energy relations [4]:

$$i\int dy \langle 0|T\{\frac{\alpha_s}{\pi}G^2(y)\frac{\alpha_s}{\pi}G^2(0)\}|0\rangle = \frac{32}{b}\langle\frac{\alpha_s}{\pi}G^2\rangle \qquad (14)$$

$$i\int dy \langle 0|T\{\bar q q(y)\frac{\alpha_s}{\pi}G^2(0)\}|0\rangle = \frac{24}{b}\langle\bar q q\rangle, \qquad (15)$$

$$b = 11-\frac{2}{3}N_f = 9$$

Suppose that the relations of the type of (14),(15) are saturated by some effective state σ of the mass $m_\sigma \simeq 0,7$ GeV. We do not insist on the existence of real narrow gluonium state with such a mass, but rather imply the effective description of 0^+ - channel. The scale 0,7 GeV is understood as that region where $\pi\pi$ -interaction becomes strong and phases vary considerably.

In the notations

$$\langle 0|\frac{\alpha_s}{\pi}G^2|\sigma\rangle = \lambda_1, \quad \langle 0|\bar q q|\sigma\rangle = \lambda_2 \qquad (16)$$

we have

$$\frac{\lambda_1^2}{m_\sigma^2} = \frac{32}{b}\langle\frac{\alpha_s}{\pi}G^2\rangle, \quad \frac{\lambda_1\lambda_2}{m_\sigma^2} = \frac{24}{b}\langle\bar q q\rangle, \quad K = \frac{\lambda_2^2}{m_\sigma^2} \qquad (17)$$

Then it follows from (12) (at $\langle\bar q q\rangle = -(0{,}25\ \mathrm{GeV})^3$, $\langle\frac{\alpha_s}{\pi}G^2\rangle = 1{,}2\cdot 10^{-2}\,\mathrm{GeV}^4$) :

$$\langle N|\bar s s|N\rangle \simeq -\frac{18}{b}\frac{\langle\bar q q\rangle}{\langle\frac{\alpha_s}{\pi}G^2\rangle}\, m \simeq 2{,}4 \qquad (18)$$

(it corresponds to $K \simeq 0{,}04\ \mathrm{GeV}^2$). This estimate does not contradict (7) ((7) would be reproduced at $K \simeq 0{,}06\ \mathrm{GeV}^2$).

We can found K also in the chiral SU(3)-limit assuming $m_s \to 0$ and keeping only the main term in K proportional to $\ln m_s^{-1}$. In this limit K is defined by the contribution of the states K^+K^-, $\eta\eta$ which are uniquely fixed by the PCAC relations:

$$\langle 0|\bar u u|K^+K^-\rangle = \frac{m_K^2}{m_s}, \quad \langle K^+K^-|\bar s s|0\rangle = \frac{m_K^2}{m_s} \qquad (19)$$

$$\langle 0|\bar u u|\eta\eta\rangle = \frac{1}{3}\frac{m_K^2}{m_s}, \quad \langle\eta\eta|\bar s s|0\rangle = \frac{4}{3}\frac{m_K^2}{m_s} \qquad (20)$$

Proceeding in the usual way [4] and saturating subtracted dispersion relation for K by K^+K^-, $\eta\eta$ states we get

$$\mathrm{Im}\, K(t)^{K^+K^-+\eta\eta} = \frac{m_K^2}{m_s}\left(1+\frac{4}{9}\right)\frac{v_K}{16\pi}, \quad v_K = \sqrt{1-\frac{4m_K^2}{t}}$$

$$K = \frac{1}{\pi}\int_{4m_K^2}^{\infty}\frac{dt}{t}\,\mathrm{Im}\,K(t) \simeq \left(\frac{m_K^2}{m_s}\right)^2\left(1+\frac{4}{9}\right)\ln\frac{M^2}{m_K^2}\cdot\frac{1}{16\pi^2} \simeq \qquad (21)$$

$$\simeq 0{,}03\ \mathrm{GeV}^2$$

For numerical estimate we put $\ln\frac{M^2}{m_K^2} = 1$ in (21). The real accuracy of SU(3) chiral estimate of K is not high. However, the value (21) is close to the previous estimate.

Finally, there is the third way to find $\langle N|\bar s s|N\rangle$ which does not appeal to sum rules at all. It implies saturation of $\langle N|\bar s s|N\rangle$ by the lowest SU(3) singlet state-gluonium σ (see (16)). Such the saturation allows one to find the ratio of nucleon matrix elements of operators $\bar s s$ and $\theta_{\mu\mu}$:

$$\frac{\langle N|\bar s s|N\rangle}{\langle N|\theta_{\mu\mu}|N\rangle} = \frac{\langle 0|\bar s s|\sigma\rangle\langle\sigma|\bar N N\rangle}{\langle 0|\theta_{\mu\mu}|\sigma\rangle\langle\sigma|\bar N N\rangle} =$$

$$= \frac{\langle 0|\bar s s|\sigma\rangle}{m_s\langle 0|\bar s s|\sigma\rangle - \frac{b}{8}\langle 0|\frac{\alpha_s}{\pi}G^2|\sigma\rangle} = \frac{\lambda_2}{m_s\lambda_2 - \frac{b}{8}\lambda_1} = \qquad (22)$$

$$= -\frac{8}{9}\frac{\langle\bar s s\rangle}{\langle\frac{\alpha_s}{\pi}G^2\rangle}\,\frac{1}{1-\frac{8}{9}\langle m_s\bar s s\rangle/\langle\frac{\alpha_s}{\pi}G^2\rangle} \simeq 1\ \mathrm{GeV}^{-1}$$

To arrive at (22) we have used the explicit form of the trace of energy-momentum tensor $\theta_{\mu\mu}$ and the relations (16), (17) for λ_1, λ_2. Further, taking into account that $\langle N|\theta_{\mu\mu}|N\rangle = m\bar{N}N$, we get from (22):

$$\langle N|\bar{s}s|N\rangle \simeq 1 \qquad (22a)$$

Thus, phenomenological properties of the scalar mesons (strong mixing $\bar{u}u, \bar{d}d \leftrightarrow \bar{s}s \leftrightarrow G^2$) and large value of matrix element $\langle N|\bar{s}s|N\rangle$ are connected and agree with each other. In spite of some difference among the numerical values, the phenomenological estimates (18),(21),(22a) and the value (7) indicate that $\langle N|\bar{s}s|N\rangle$ indeed exceeds unity or at least is comparable to it.

4. S-quark and the nucleon mass

In this section the effect of s-quark on the nucleon mass is shown to be large.

To begin with, note that K enters not only the expression (12) but also determines the variation of $\langle \bar{u}u\rangle$ with s-quark mass (to compare with [4])

$$\frac{d}{dm_s}\langle \bar{u}u\rangle = -i\int dy \langle 0|T\{\bar{u}u(0)\,\bar{s}s(y)\}|0\rangle = \qquad (23)$$
$$= -K = -(0{,}04 \div 0{,}06)\,\mathrm{GeV}^2$$

Then

$$\left|\frac{\langle \bar{u}u\rangle_{m_s=140\,\mathrm{MeV}} - \langle \bar{u}u\rangle_{m_s=0}}{\langle \bar{u}u\rangle_{m_s=140\,\mathrm{MeV}}}\right| \sim 0{,}4 \div 0{,}6 \qquad (23a)$$

Considerable decrease of $|\langle \bar{u}u\rangle|$ with the variation of m_s from 0,14 GeV to $m_s = 0$ does not seem very surprising since another condensates, e.g. $\langle G^2\rangle$ [4] possess the analogous properties. From the microscopic point of view the decrease of absolute values of vacuum matrix elements with the decrease of the s-quark mass seems quite natural since topologically-nontrivial vacuum fluctuations, e.g. instantons, are supressed by light quarks [8,9].

Consider now what would happen with the nucleon mass if m_s were zero. As it is known, the nucleon mass is determined by the trace of energy-momentum tensor $\theta_{\mu\mu}$ and in the limit $m_u = m_d = 0$ it is equal to

$$m = m_s\langle N|\bar{s}s|N\rangle - \frac{b}{8}\langle N|\frac{\alpha_s}{\pi}G^2|N\rangle \qquad (24)$$

Adopting the values (7),(18),(22a) for $\langle N|\bar{s}s|N\rangle$ we conclude that considerable part of the nucleon mass (from 15% for (22a) to 50% for (7)) is due to the strange quark. How does this fact agree with the SU(3) relations for baryon masses which experimentally hold with high accuracy? The answer is that $\langle N|\bar{s}s|N\rangle$ is determined by the graph of fig. 1 being SU(3) singlet under the assumption of σ-dominance in the correlators of the type (9). The validity of this statement is supported by the low-energy theorems (15): their RHS for strange and nonstrange flavours as well as vacuum expectation values $\langle \bar{d}d\rangle$, $\langle \bar{s}s\rangle$ differs by no more than 20%.

Therefore the considered contribution is an effective addition to SU(3) singlet G^2-term in (24).

As it is seen from (24) m can vary by (15 ÷ 50)% with the decrease of m_s to zero. How one can understand this fact from the point of view of QCD sum rules? Remind that in the nucleon mass sum rules the information on m is contained in the vacuum condensates $\langle \bar{q}q\rangle, \langle G^2\rangle, \ldots$ varying with m_s considerably (see [4,9] and also (23)). It is important that this variation certainly proceeds in the necessary direction: absolute values of condensates decrease with the decreasing m_s. This leads to smaller scale in the sum rules and finally, to decrease of m. It is difficult to make any quantitative conclusions since the variation of the scale in the sum rules for $\Pi(p^2)$ (see (11)) with the variation of m_s is determined by a number of factors such as the variation of higher states contribution, the change of the ratio $\langle \bar{q}Gq\rangle / \langle \bar{q}q\rangle$ and so on.

4. Neutron electric dipole moment. Matrix element $\langle NIG\tilde{G}N\rangle$

As an application of the method described we would like to point out the essential change of the estimates for neutron dipole moment in the models where CP-violation arises in the Higgs sector. The reason is that the Higgs boson-quark couplings are proportional to the quark masses. Therefore CP-odd effects for S -quarks are enhanced by the factor $m_s/m_{u,d}$. Besides, additional large factor $ctg^2\theta_c$ (θ_c is the Cabibbo angle) can arise here. Due to the large nucleon matrix elements of s -quarks it leads to the enhancement of neutron electric dipole moment (EDM).

Consider for example, neutron EDM in the Weinberg model of CP-violation. It was shown in ref. [10] that it is mainly due to the strange quarks described by the diagrams 4a,b. Theoretically, these diagrams are singled out by their logarithmic divergency in the SU(3)xSU(3) symmetry limit. The crosses in figs. 4a,b mark CP-odd $KN\Sigma$ vertex caused by the operator

$$H_s = \Lambda g\,\bar{s}\sigma_{\mu\nu}\gamma_5\frac{\lambda^n}{2}G^n_{\mu\nu}s \tag{25}$$

This interaction arises due to the diagram 5. The constant Λ numerically equals to

$$\Lambda \simeq -1{,}5\cdot 10^{-23}\,sm \tag{26}$$

By means of PCAC hypothesis CP-odd constant is transformed to the form

$$g'_{KN\Sigma} = \langle K\Sigma^-|H_s|n\rangle = \frac{\Lambda}{f_K}\langle\Sigma^-|ig\,\bar{s}\sigma_{\mu\nu}\frac{\lambda^n}{2}G^n_{\mu\nu}u|n\rangle \tag{27}$$

The resulting matrix element was estimated in the nonrelativistic quark model to be

$$\langle\Sigma^-|ig\bar{s}\sigma_{\mu\nu}\frac{\lambda^n}{2}G^n_{\mu\nu}u|n\rangle \simeq 0{,}12\,GeV^2 \tag{28}$$

The strange quark contribution to neutron EDM was evaluated in this way to be [10]

$$D_n(s) \simeq -8\cdot 10^{-25}\,e\cdot sm \tag{29}$$

In fact, the number (28) seems to be strongly underestimated. Really, we have from (13a):

$$\langle\Sigma^-|\bar{s}\hat{O}u|n\rangle \simeq \frac{m\,iK_{us}}{-\langle\bar{q}q\rangle},\quad K_{us} = \int dy\langle 0|T\{\bar{s}\hat{O}u(y)\cdot\bar{u}s(0)\}|0\rangle \tag{30}$$

The correlator arising here can be found in the following way:

$$\begin{aligned} 0 &= i\partial^\mu_x\int dy\langle 0|T\{\bar{u}\gamma_\mu s(x)\bar{s}\hat{O}u(y)\}|0\rangle = \\ &= (m_s-m_u)K_{\bar{u}s} + \langle 0|[u^+s,\bar{s}\hat{O}u]|0\rangle = \\ &= (m_s-m_u)K_{\bar{u}s} + \langle 0|(\bar{u}\hat{O}u - \bar{s}\hat{O}s)|0\rangle \end{aligned} \tag{31}$$

which is equivalent to method of ref. [11].

As a result, we find

$$\langle\Sigma^-|\bar{s}\hat{O}u|n\rangle \simeq f_g m_0^2\cdot\frac{m}{m_s-m_u} \simeq -1\,GeV^2 \tag{32}$$

Here $m_0^2 \equiv \langle\bar{q}ig\sigma_{\mu\nu}\frac{\lambda^n}{2}G^n_{\mu\nu}q\rangle\langle\bar{q}q\rangle^{-1} \simeq 1\,GeV^2$ [6], $f_g = \langle(\bar{s}\sigma Gs - \bar{u}\sigma Gu)\rangle\langle\bar{u}\sigma Gu\rangle^{-1}$ To determine f_g we use the method suggested in ref. [12] to calculate $f = \langle(\bar{s}s-\bar{u}u)\rangle\langle\bar{u}u\rangle^{-1} \simeq -0{,}23$. In this way we get

$$f_g = -\frac{2}{3}f \simeq -0{,}15 \tag{33}$$

(see the Appendix). Note that there could be other contributions comparable to (30) (see, e.g., fig. 6). Explicit calculation of them leads to zero result for the specific choice of the baryon currents in the form $\varepsilon^{abc} q_1^a \gamma^\mu (q_2^b C \gamma_\mu q_3^c)$) [7].

The result (32) exceeds by an order of magnitude the best experimental limit on the neutron EDM [13]. Thus, the Weinberg model of CP-violation is definitely excluded by the experimental data.

Earlier the same conclusion was made in the work [14] where the contribution of the diagram with the neutral Higgs boson to the neutron EDM was estimated to be at the level $\sim 10^{-22}$ e·sm under the assumption that the coupling constants of the neutral Higgs boson are of the same order of magnitude as those of the charged bosons. It seems to us, however, that taking into account the formfactors in the Higgs vertices leads to the decrease of this number down to $\sim 10^{-23}$ e·sm. Besides, there is a more essential circumstance [14,15]. We have no information at all on the CP-violation in the system of neutral Higgs bosons. Therefore the result of ref. [14] can be interpreted merely as a limit on the CP-odd constants of these particles in the considered model.

Among other interesting applications of the method considered we would like to point out the estimate of the matrix element $\langle N | G^n_{\mu\nu} \tilde{G}^n_{\mu\nu} | N \rangle$ which determines the interaction of a pseudoscalar Higgs boson with nucleon. Using low-energy theorem [16],

$$\int dy \langle 0 | T\{ \frac{\alpha_s}{8\pi} G\tilde{G}(y) \bar{q}\gamma_5 q(0) \} | 0 \rangle = -\frac{1}{N_f} \langle \bar{q} q \rangle \qquad (35)$$

(in our notations) and formula (13b) we find

$$\langle N | N_f \frac{\alpha_s}{4\pi} G\tilde{G} | N \rangle \simeq 2m \bar{N} i \gamma_5 N \qquad (36)$$

In the chiral SU(3) limit this formula corresponds to the value $g_A^S = 1$ for the constant g_A^S of the axial SU(3)-scalar nucleon current:

$$\langle N | \sum_{q=u,d,s} \bar{q} \gamma_\mu \gamma_5 q | N \rangle \simeq \bar{N} \gamma_\mu \gamma_5 N \qquad (37)$$

One can easily get this result by means of the expression for anomaly: $\sum_q \partial^\mu \bar{q} \gamma_\mu \gamma_5 q = \frac{\alpha_s}{4\pi} G\tilde{G}$ The SU(6)-symmetry quark model also gives $g_A^S = 1$. Other estimates lead to the values of g_A^S somewhat smaller than unity [17,18].

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Appendix

Here the method of calculation of the value f_g (see (33)) is described. In ref. [12] the correlators

$$i\int dx e^{iqx} \langle 0|T\{j(x)\bar{s}\gamma_5 u(0)\}|0\rangle \qquad (1A)$$

$$i\int dx e^{iqx} \langle 0|T\{j(x)\bar{s}\gamma_\lambda\gamma_5 u(0)\}|0\rangle \qquad (2A)$$

are considered, where $j = \bar{u}\gamma_\mu\gamma_5 i\overleftrightarrow{D}_\nu s$. Keeping only K - -meson contribution in (1A),(2A) the sum rules for them of the form

$$\exp\left(-\frac{m_K^2}{M^2}\right) f_K^2 \frac{m_K^2}{m_s+m_u} C \simeq \frac{m_s-m_u}{8\pi^2} M^2 + \frac{\langle(\bar{s}s-\bar{u}u)\rangle}{1} + \ldots \qquad (3A)$$

$$\exp\left(-\frac{m_K^2}{M^2}\right) f_K^2 C \simeq \frac{m_s^2-m_u^2}{4\pi^2} + \frac{\langle(m_u\bar{u}u - m_s\bar{s}s)\rangle}{M^2} + \ldots \qquad (4A)$$

are obtained, where C is a constant. Then

$$\frac{\langle\bar{u}u\rangle}{\langle\bar{s}s\rangle} \approx 1 + \frac{m_K^2}{M^2}\,\frac{m_s-m_u}{m_s+m_u} \qquad (5A)$$

where $M = m_{K^*}$ is the typical scale in the sum rules for 0^-, 1^+ channels.

To estimate f_g let us choose the current j in the form $\bar{u}\gamma_\mu\gamma_5\frac{\lambda^n}{2}G^n_{\mu\nu}s$. As a result, we get the following sum rules:

$$\exp\left(-\frac{m_K^2}{M^2}\right) f_K^2 \frac{m_K^2}{m_s+m_u} C' \simeq \frac{\langle(\bar{s}\sigma G s - \bar{u}\sigma G u)\rangle}{1} + \ldots \qquad (6A)$$

$$\exp\left(-\frac{m_K^2}{M^2}\right) f_K^2 C' \simeq \frac{2}{3}\,\frac{\langle(m_u\bar{u}\sigma G u - m_s\bar{s}\sigma G s)\rangle}{M^2} + \ldots \qquad (7A)$$

They lead to

$$\frac{\langle\bar{u}\sigma G u\rangle}{\langle\bar{s}\sigma G s\rangle} \approx 1 + \frac{2}{3}\,\frac{m_K^2}{M^2}\,\frac{m_s-m_u}{m_s+m_u}. \qquad (8A)$$

and we arrive at the estimate (33).

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Figure captions

Fig. 1. Contribution to the correlator (8) from the expectation value (9).

Fig. 2. Contribution to the polarization operator (11) from the expectation value $\langle \bar{q}q \rangle$. Grosses denote the vacuum fields.

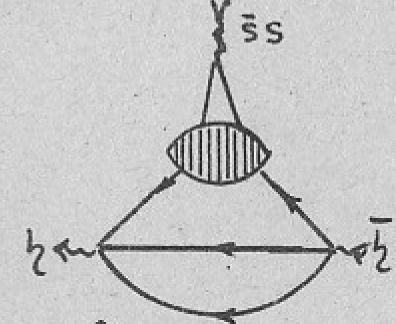
Fig. 3. An example of the contribution of $\bar{u}u$ at small distances to $\langle N|\bar{u}u|N \rangle$ (asymptotic loop).

Fig. 4. Diagrams for neutron EDM in the Weinberg model of CP-violation.

Fig. 5. The diagram leading to the operator $\bar{s}\sigma G\gamma_5 s$ in the effective CP-odd Hamiltonian (H is the Higgs boson, g is the gluon).

Fig. 6. An example of the contribution of $\bar{s}\sigma G u$ at small distances to $\langle \Sigma^- | s\sigma G u | n \rangle$ (contribution from the VEV $\langle q\bar{q}q\bar{q}G \rangle$).

Fig. 7. The diagram for the neutron EDM due to the neutral Higgs boson exchange.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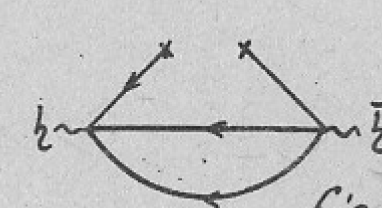
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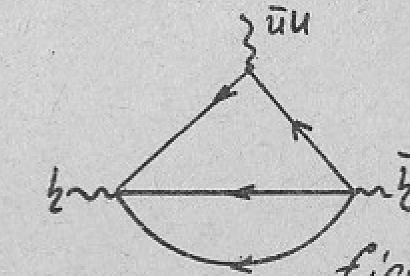
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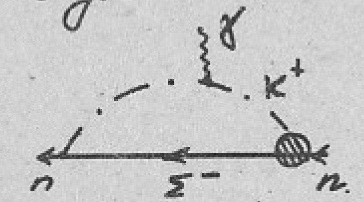
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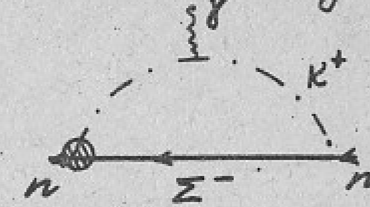
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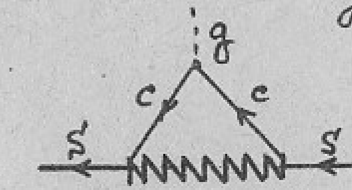
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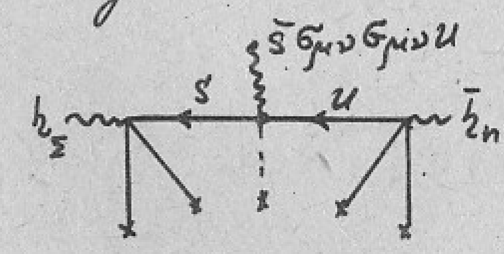
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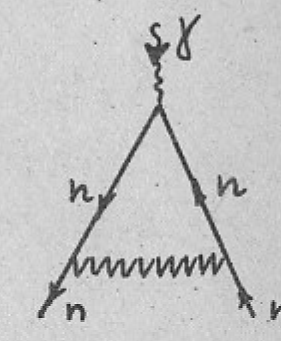
fig.7

В.М.Хацимовский, И.Б.Хриплович,

А.Р.Житницкий

СТРАННЫЕ КВАРКИ В НУКЛОНЕ И ДИПОЛЬНЫЙ МОМЕНТ НЕЙТРОНА

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