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SELF-CONSISTENT CALCULATION OF THE WEAK CONSTANTS IN THE PARITY NONCONSERVING NUCLEAR FORCES.

Effective PNC Hamiltonian in $SU(2)_L \times U(1) \times SU(3)_c$. PNC in the πNN Vertex

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1. INTRODUCTION

This work is devoted to the consequences of the standard model $SU(2)_L \times U(1) \times SU(3)_c$ /1-3/ for the parity nonconserving (PNC) nucleon-nucleon (NN) interactions. We carry out the analysis in the framework of the usually accepted representation of the PNC NN forces through the π-, ρ-, ω- meson exchanges with the parity nonconservation within one of the MNN ($M = \pi, \rho, \omega$) vertices (see, e.g., reviews/4,5/, papers/6,7/). Here, the main theoretical problem is the calculation of the PNC MNN vertices defined by the phenomenological Hamiltonian x)

$$\mathcal{H}^{PNC}_{MNN} = \frac{1}{\sqrt{2}} h_\pi \bar{N} (\vec{\tau} \times \vec{\pi})^3 N +$$

$$+ \bar{N} \left(h^0_\rho \vec{\tau} \vec{\rho}_\mu + h^1_\rho \rho^3_\mu + h^2_\rho \frac{3\tau^3 \rho^3_\mu - \vec{\tau} \vec{\rho}_\mu}{2\sqrt{6}} \right) \gamma^\mu \gamma_5 N + \tag{1}$$

$$+ N \left(h^0_\omega \omega_\mu + h^1_\omega \tau^3 \omega_\mu \right) \gamma^\mu \gamma_5 N .$$

An activity in this field is necessary at least by the following reasons.

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x) We use the Bjorken-Drell metric with $\gamma_5 = \begin{pmatrix} 0 & -1 \\ -1 & 0 \end{pmatrix}$, our h_π is the same as f_π in refs./5,7/. We neglect the PNC ρNN vertex with derivatives ($\propto h^{1'}_\rho$), because it plays a negligible role in the PNC NN interactions (see ref./8/).

First, the PNC NN interactions give the unique possibilities for the study of all the components of the standard model, which at present is the low-energy benchmark for the most grand unified models. Indeed, unlike the $\Delta S = 1$ or/and $\Delta C = 1$ nonleptonic decay amplitudes, the PNC MNN vertices are generated not only by the charged but also neutral hadron current interactions, and therefore, bear information on unification scheme of all three types of the fundamental interactions: weak, electromagnetic and strong. So, the $\Delta I = 1$ PNC NN transitions are determined almost completely by the neutral current interactions, and the modification of the bare (i.e., determined by the weak Lagrangian) PNC quark amplitudes by quark-gluon interactions may tell on the values of h_M very strongly (up to an order of magnitude and sign).

Second, such an analysis may have interesting consequences for some extensions of the standard model (e.g., for the supersymmetric ones, see ref. /9/).

A consistent consideration of our subject however encounters essential difficulties because there is no yet an adequate technique for the transition from the quark interaction Lagrangian of the standard model to the hadronic amplitudes. The available calculations of the constants h_M necessarily include dissimilar tricks and approximations, that is the main origin of large uncertainties of the results of these calculations (see, e.g., refs. /5,10/.

In the paper /7/ Desplanques, Donoghue and Holstein have carried out the unified treatment of h_M based upon $SU(6)_W$ arguments supplemented by quark technique and have analysed some of these uncertainties.

As a result, they estimated the intervals of possible values of h_M in the standard model and within the intervals found the so-called "best values" of h_M $\left(h_M^{b.v.}\right)$, which at present are widely used for the estimations of the PNC effects in the standard model. The "best values", however, are actually semiphenomenological ones, because to obtain them, arbitrary (fitting) parameters have been used *). This circumstance hinders the use of $h^{b.v.}$ to conclude on consequences of the standard model itself.

*) Moreover some of these parameters have rather an unnatural interpretation: large contribution of a quark sea to some h_M; a strong breaking of the $SU(6)_W$ symmetry for nonfactorizable parts of h_M.

For this aim a self-consistent parameter-free calculational framework for h_M is required. This framework also has to enable one to sift possible artifacts of calculational methods included in it.

Our realization of this plan is close in the form to the unified treatment of h_M of ref. /7/, although differs from it in the basic points. We begin with a treatment of the effective Hamiltonian of PNC hadron-hadron interactions in the standard model, which is the local operator in the first order generating the PNC MNN vertices: $h_M \propto \langle MN' | \mathcal{H}^{PNC} | N \rangle$. The essential point here is the choice of a low renormalization scale of the operators of the Hamiltonian (μ is near the characteristic inverse confinement radius). This enables us to consider hadronic matrix elements of this operators taking into account only valence quarks. Then we expound the overall structure of the PNC MNN vertices determined by this effective Hamiltonian. We pay special attention to a reconstruction of quark mechanisms responsible for different contributions to h_M. After that we proceed to the calculation of h_M. Here, the main difficulty is the calculation of nonfactorizable (NF) (i.e., nonreducible to products of hadron currents) contributions to h_ρ and h_ω. We solve the problem proceeding from the approximate SU(6) symmetry of the NF parts of the PNC MNN vertices and calculate $h^{NF}_{\rho,\omega}$ in the MIT model /11/ (like h^{NF}_π). Our values of h_M are within the intervals determined in ref. /7/, moreover $h_\pi \simeq \frac{1}{3} h_\pi^{b.v.}$, and $h_{\rho,\omega}$ agree with $h^{b.v.}_{\rho,\omega}$. We analyse also the partial contributions to h_M of the charged and neutral currents and a role of the quark-gluon interactions in the formation of h_M. The concluding part of the work is devoted to the experimental consequences of our results. We observe a rather good agreement of our values of h_M with the experimental data of the whole and concentrate our attention on the processes valuable for more elaborate conclusions. Our work is organized in the form of three issues. In the remaining part of this issue we consider the effective Hamiltonian $\mathcal{H}^{PNC}$ (sect.2) and calculate h_π (sect.3). The second issue (II) contains the calculation of $h_{\rho,\omega}$; in the third issue (III) we discuss the experimental consequences of our results.

2. THE EFFECTIVE HAMILTONIAN OF THE PNC HADRON-HADRON INTERACTIONS IN THE STANDARD MODEL

The general structure of the effective Hamiltonian of weak hadron-hadron interactions in the unified gauge theories has been studied in papers [12,13]. Their main result is that a part of $\mathcal{H}$, that is not reduced to the renorlamization of the starting Lagrangian, is finite, its leading terms being of the order $O(G \sim g_w^2/M_w^2)$ (g_w is the semiweak coupling constant). The explicit form of $\mathcal{H}_{\Delta S=1}$ in the standard model with massless quarks in the second order in g_w and in all orders in the quark-gluon coupling constant g_s (in the leading log approximation) has first been found in ref. [14]. There have been observed the dynamical, owing to the quark-gluon interactions, enhancement of octet transitions and suppression of 27-plet quark transitions. In ref. [15] analogous effects have been found for PNC Hamiltonian $\mathcal{H}^{PNC}_{\Delta S=0}$. The next step in the discovery of the structure of $\mathcal{H}$ in the standard model was to take into account the difference of the quark mass scale: $m_c \gg m_{u,d,s}$. So, in ref. [16] it has been shown that the elimination of the logarithmic contributions of c - quarks from the operator part of $\mathcal{H}_{\Delta S=1}$ leads to the appearance in $\mathcal{H}_{\Delta S=1}$ of new, including the neutral right-handed currents, operator structures which belong to the octet representation (the so-called "penguins"). In the papers [17,18] analogous effects have been studied for the $\Delta I=1$ part of $\mathcal{H}^{PNC}_{\Delta S=0}$, and in refs. [19,20] for all the isotopic parts of $\mathcal{H}^{PNC}_{\Delta S=0}$. We are interested in the effective Hamiltonian $\mathcal{H}^{PNC}$ which determines the low-energy (E $\ll$ 1 GeV) PNC interactions of light hadrons. Here, following the calculation scheme of ref. [20], we shall give $\mathcal{H}^{PNC}$ in its final form and dwell only upon its features most important for the following.

Within the standard model *) in the second order in g_w and in all orders of the leading logs in the quark-gluon interactions the Hamiltonian $\mathcal{H}^{PNC}$ has the form:

$$\mathcal{H}^{PNC} = \mathcal{H}^{PNC}_{\Delta S=1} + \mathcal{H}^{PNC}_{\Delta S=0}, \tag{2}$$

$$\mathcal{H}^{PNC}_{\Delta S=1} = \sqrt{2}\, G\, sc \left\{ -\frac{1}{10} K_w^{d_{84}} K^{d_{27}} O^{27}_{3/2,\,1/2} + \right.$$

*) For simplicity we confine ourselves to the GIM sector [2] (the number of flavours $n_f = 4$); subsequent results are practically not sensitive to the expansion of the quark basis to $n_f = 6$ (see ref. [21]).

$$\left. + \left(\frac{1}{20} [1\ 1\ 1\ 1]\, K_w^{d_{84}} + \frac{1}{4} [-1\ 1\ -1\ 1]\, K_w^{d_{20}} \right) K^{d_8} O^8_{1/2} \right\}, \tag{3}$$

$$\mathcal{H}^{PNC}_{\Delta S=0} = \sqrt{2}\, G \left\{ K_w^{d_{84}} K^{d_{27}} \left[\frac{1}{6} c^2 O^{27}_2 - \frac{1}{10} s^2 O^{27}_1 + \frac{1}{120} (1-4s^2) O^{27}_0 \right] + \right.$$

$$+ \left(\frac{1}{20} [1\ 1\ 1\ 1]\, K_w^{d_{84}} + \frac{1}{4} [-1\ 1\ -1\ 1]\, K_w^{d_{20}} \right) K^{d_8} \left[s^2 O^8_1 + \frac{1}{\sqrt{3}} (2-3s^2) O^8_0 \right] +$$

$$\left. + \left(-\frac{1}{120} [1\ 1]\, K_w^{d_{84}} + \frac{1}{30} [-1\ 4]\, K_w^{d_1(4)} \right) K^{d_1(3)} O^1_0 \right\} +$$

$$+ \sqrt{2}\, G (1-2s_w^2) \left\{ K_Z^{d_{84}} K^{d_{27}} \left(-\frac{1}{6} O^{27}_2 + \frac{1}{10} O^{27}_1 + \frac{1}{60} O^{27}_0 \right) + \right.$$

$$+ \left(-\frac{1}{20} [1\ 1\ 1\ 1]\, K_Z^{d_{84}} + \frac{1}{4} [-1\ 1\ -1\ 1]\, K_Z^{d_{20}} \right) K^{d_8} \left(O^8_1 + \frac{1}{\sqrt{3}} O^8_0 \right) +$$

$$\left. + \left(\frac{1}{60} [1\ 1]\, K_Z^{d_{84}} + \frac{1}{60} [-1\ 4]\, K_Z^{d_1(4)} \right) K^{d_1(3)} O^1_0 \right\} -$$

$$- \sqrt{2}\, G \frac{1}{3} s_w^2 \, [1\ 0\ 0\ 0]\, K_Z^{d_{15}} \left\{ K^{d_8} \left(O^8_1 + \frac{1}{\sqrt{3}} O^8_0 \right) - \right.$$

$$\left. - \frac{1}{6} \begin{pmatrix} K^{d_1(3)} O^1_0 \\ K^{d_1(3)} O^1_0 \end{pmatrix} \right\}. \tag{4}$$

The Hamiltonian is represented by the linear combination of the local four-quark operators O^R_i renormalized at the point $\mu \simeq \mu_0$, where μ_0 is a parameter of the infrared cut-off of the quark and gluon loop momenta. This cut off is introduced at the characteristic inverse radius of the confinement and takes phenomenologically into account the colour confinement. In our case $\mu_0 \simeq 0.2$ GeV. The upper indices of the operators denote the $SU(3)$ representations which they belong to, and the lower ones show their properties in the isospin space (ΔI). The explicit form of operators O^{27}, is given in Appendix; O^{8} and O^{1} are defined as follows:

$$O_i^8 = \begin{pmatrix} O(\Lambda_i, 1) \\ O^c(\Lambda_i, 1) \\ O(1, \Lambda_i) \\ O^c(1, \Lambda_i) \end{pmatrix}, \qquad O_0^1 = \begin{pmatrix} O(1,1) \\ O^c(1,1) \end{pmatrix}, \qquad (5)$$

where $O(M,N) = :\bar{q}_i \gamma_\mu \gamma_5 M q_i \bar{q}_j \gamma^\mu N q_j:$, $O^c(M,N) = :\bar{q}_i \gamma_\mu \gamma_5 M q_j \bar{q}_j \gamma^\mu N q_i:$, $q_i \equiv q_i(0)$, a sum over colour indices i,j = 1, 2, 3 is assumed; M, N are 3 x 3 matrices in the flavour space; $\Lambda_{1/2} = \lambda_6$, $\Lambda_1 = \frac{1}{2}\lambda_3$, $\Lambda_0 = \frac{1}{2}\lambda_8$ (λ_i are the Gell-Mann matrices).

All the parameters of the standard model: $s \equiv \sin\theta_c$ $(c \equiv \cos\theta_c)$, $s_w \equiv \sin\theta_w$, g_w, g_s are contained in the coefficient functions of O^R. Summation of the leading log corrections (here-through the renormalization group equations in the massless scheme of renormalization[22]) leads to the matrices K^d unequal to identity (as $g_s \to 0$, $K^d \to 1$). The numbers K are defined by $K_{w,z} = \alpha_4(m_c^*)/\alpha_4(M_{w,z})$, $K = \alpha_3(\mu_0)/\alpha_3(m_c^*)$, where $\alpha_n(\mu) = 12\pi/[(33-2n)\ln(\mu^2/\Lambda_n^2)]$ is the effective strong interaction coupling $(g_s^2(\mu)/4\pi)$ in the theory with n flavours; m_c^* is the mass of a charmed quark with a special choice of the renormalization point, namely with μ, which is a solution of the equation $m_c(\mu) = \mu$ *). The matrices d_R (proportional to the matrices of the anomalous dimensions of the operators O^R) have the form

$$d_{84} = -6/25, \quad d_{27} = -6/27, \quad d_{20} = 12/25; \qquad (6)$$

*) Hence $m_c^* = m_c(m_c^*)$.

$$d_{15} = d_{AD}(4), \quad d_8 = d_{AD}(3),$$

$$d_{AD}(n) = -\frac{1}{33-2n}\begin{pmatrix} -2/3 & 2 & -3 & 9 \\ \frac{27-4n}{6} & -\frac{27-4n}{2} & 9/2 & 21/2 \\ -11/3 & 11 & 0 & 0 \\ 9/2 & 21/2 & 9/2 & -27/2 \end{pmatrix}; \qquad (7)$$

$$d_1(n) = -\frac{1}{33-2n}\begin{pmatrix} -11/3 & 11 \\ \frac{27-2n}{3} & 2n-3 \end{pmatrix}. \qquad (8)$$

To find numerical values of the coefficient functions of $\mathcal{H}^{PNC}$, we should choose the values of the parameters Λ_4, Λ_3, m_c^*. Note, only two of these three parameters are independent. Indeed, the matching condition for the amplitudes $\langle u,d,s \,|\, O_4(u,d,s,c) \,|\, u,d,s\rangle$ and $\langle u,d,s \,|\, O_3(u,d,s) \,|\, u,d,s\rangle$ in the log approximation at $\mu = m_c^*$ leads to $\alpha_3(m_c^*) = \alpha_4(m_c^*)$; hence, $\Lambda_3 = m_c^*(\Lambda_4/m_c^*)^{25/27}$ (see also [23]). We use Λ_4 = 80 MeV that complies with the charmonium data [24], and m_c^* = 1.27 GeV that follows from the analysis of the e^+e^- - annihilation experimental data through the QCD sum rules (see, e.g., review[25] and references therein). Hence, $\Lambda_3 \simeq$ 100 MeV and at μ_0 = 0.2 GeV we obtain $K_w \simeq K_z \simeq 2.5$ and $K \simeq 3.6$.

It is convenient to write the final expression for $\mathcal{H}^{PNC}$ in terms of O^{27} (see Appendix) and of the operators:

$$\begin{pmatrix} O_i^S \\ O_i^A \\ O_i^5 \\ O_i^6 \end{pmatrix} = \begin{pmatrix} 1 & 1 & 1 & 1 \\ -1 & 1 & -1 & 1 \\ 0 & 1 & 0 & -1 \\ 1 & 0 & -1 & 0 \end{pmatrix} O_i^8, \qquad \begin{pmatrix} O^{1S} \\ O^{1A} \end{pmatrix} = \begin{pmatrix} 1 & 1 \\ -1 & 1 \end{pmatrix} O_0^1. \qquad (9)$$

The explicit form of the operators is given in Appendix, too. It is easy to see that the operators O^{27} , O^S and O^{1S} (O^A , O^{1A}) are symmetric (antisymmetric) with respect to the permutation of the colour indices of the second and fourth or first and third quarks. The operators O^5 and O^6 have no definite colour symmetry and have the left-right ("penguin") structure: $O_{LR} \propto O(M,N) - O(N,M)$.

As a result we have:

$$\mathcal{H}^{PNC}_{\Delta S=1} = \sqrt{2}\, G \Big(c^{27}_{\Delta S=1}\, O^{27}_{3/2,\,1/2} + c^S_{1/2}\, O^S_{1/2} + c^A_{1/2}\, O^A_{1/2} + c^5_{1/2}\, O^5_{1/2} + c^6_{1/2}\, O^6_{1/2} \Big), \qquad (10)$$

$$\mathcal{H}^{PNC}_{\Delta S=0} = \sqrt{2}\, G \Big[c^{27}_2\, O^{27}_2 + \sum_{i=0,1} \big(c^{27}_i O^{27}_i + c^S_i O^S_i + c^A_i O^A_i + c^5_i O^5_i + c^6_i O^6_i \big) + c^{1S} O^{1S} + c^{1A} O^{1A} \Big]. \qquad (11)$$

The numerical values of the coefficient functions C^R_i at s = 0.23 and S^2_W = 0.23 and the obtained values of K are given in Tables 1, 2 . For comparison we give also the bare values of C^R_i $(K_{W,Z} = K = 1)$. It is seen, the quark-gluon interactions result in the increase (in absolute value) in the coefficients of the antisymmetric operators O^A and O^{1A} and in the decrease in the coefficients of the symmetric operators O^{27} , O^S and O^{1S} . The characteristic ratios $|C(g_S \neq 0)/C(g_S = 0)|$ are ~ 2.5 and ~ 0.5, respectively. It is to be remembered however that we have extrapolated the coefficient functions found in the leading log approximation valid at $\alpha_S(\mu) \ll 1$, into the region $\mu \gg m_c$, where $\alpha_S(\mu) \sim 1$, and therefore, essential contributions may come from nonleading log terms $\sim g_S^{2n} (\ln m/\mu_0)^{2m}$ $(n > m)$ and from nonlogarithmic terms. All this may produce an especially strong influence on the behaviour of $C^5_{1/2}$ and $C^6_{1/2}$ which gather their values mainly at $\mu > m_c$. So, it is reasonable to assume that the differences $C(g_S \neq 0) - C(g_S = 0)$ are reproduced by the leading logs up to a factor of $\sim 1.5 \div 2$ /26/.

It is important to note here that the choice of the so low renormalization point $(\mu \simeq \mu_0)$ is caused by the following circumstan-

<u>Table 1.</u> Coefficient functions of $\mathcal{H}^{PNC}_{\Delta S=1}$ (10). (s = 0.23) In brackets their values at α_S = 0 are given.

$C^{27}_{\Delta S=1}$	$C^S_{1/2}$	$C^A_{1/2}$	$C^5_{1/2}$	$C^6_{1/2}$
−0.014	0.0049	0.15	−0.011	0.0020
(−0.022)	(0.011)	(0.056)	(0)	(0)

<u>Table 2.</u> Coefficient functions of $\mathcal{H}^{PNC}_{\Delta S=0}$ (11). (s = 0.23, S^2_W = 0.23)

C^{27}_2	C^{27}_1	C^{27}_0	C^S_1	C^A_1	C^5_1	C^6_1
0.041	0.029	0.0094	−0.027	0.46	−0.080	−0.029
(0.068)	(0.049)	(0.016)	(−0.044)	(0.17)	(0)	(−0.038)

C^S_0	C^A_0	C^5_0	C^6_0	C^{1S}	C^{1A}
0.0072	0.95	−0.098	−0.0078	0.016	0.17
(0.026)	(0.35)	(0)	(−0.022)	(0.071)	(0.099)

ces very essential for further considerations: the point μ separates large $(\mu < p < M_W)$ and small $(\mu_0 < p < \mu)$ virtual momenta between the coefficient functions and the operator part of the effective Hamiltonian; at $\mu = \mu_0$ the log contributions of all the loops, including those of the quark sea, are completely concentrated at the coefficient functions, and therefore, the hadronic matrix elements of the operators $O(\mu=\mu_0)$ in the log approximation are determined by only valence quarks*) (see ref. /27/).

3. PARITY NONCONSERVATION IN THE πNN VERTEX

The PNC πNN vertex (see (1)) plays a distinct role in the description of PNC NN interactions: On the one hand h_π determines the one-pion PNC exchange intensity, i.e., the long-range part ($r_\pi \simeq$ 3 fm) of PNC NN forces, on the other hand the value of h_π is intimately connected with the structure of neutral hadronic currents and with the magnitude of such fundamental parameters of theory as the quark vacuum condensate density $\langle 0|\bar{q}q|0\rangle$ and the quark masses m_q.

At the same time the recent estimates of h_π in the standard model are scattered in a very large interval: $2.5 \times 10^{-7} \lesssim h_\pi \lesssim 11 \times 10^{-7}$ (see papers /7,17,18,28/ and reviews /5,29/). We may distinguish the following main origins of this situation: a different choice of the renormalozation point μ in the $\mathcal{H}^{PNC}$; the use of different values of m_q in the calculation of the factorizable contributions to h_π; dissimilarity of the calculational methods for the nonfactorizable part of h_π. In more detail we dwell upon these points in the main text.

We begin our consideration of the PNC πNN vertex with clarification of the overall structure of the $\pi B'B$ (B, B' are nucleons and/or hyperons) vertices determined by the matrix elements

$$\langle \pi^a B^b | \mathcal{H}^{PNC} | B^c \rangle = A(B^c_a) \frac{i}{\sqrt{(2\pi)^3\, 2K^0}}\, \bar{u}^b u^c \tag{12}$$

(indices a, b, c identify the components of the pion isotriplet and the nucleon octet; in this notation $h_\pi = A(n^0_-)$).

3.1. Overall Structure of the PNC πB'B Vertices

Let us write down the effective Hamiltonian $\mathcal{H}^{PNC}$ in the form

$$\mathcal{H}^{PNC} = \sqrt{2}\, G \sum_{M,N} C^{MN} O^{MN}, \tag{13}$$

*) The choice $\mu=\mu_0$ is possible owing to the small value of $\Lambda_3 \simeq 0.1$ GeV, then $\alpha_3(\mu_0) \simeq 1$.

where $O^{MN} = :\bar{q}Mq\bar{q}Nq:$, M, N are matrices in the spinor x flavour x colour space, and consider the partial amplitude $A^{MN} = \sqrt{(2\pi)^3 2K^0}\, \langle \pi^a B^b | O^{MN} | B^c \rangle$.

The standard reduction technique and PCAC representation for the interpolating pion field

$$\pi^a(x) = -\frac{\sqrt{2}}{f_\pi m_\pi^2}\, \partial^\mu \bar{q}(x) P^a_\mu q(x), \qquad P^a_\mu = \gamma_\mu \gamma_5 \frac{\tau^a}{2}, \tag{14}$$

($f_\pi \simeq$ 132 MeV is the decay constant of $\pi \to \mu\nu$) allow us to write down

$$A^{MN} = -\frac{i\sqrt{2}\,(-\kappa^2 + m_\pi^2)}{f_\pi m_\pi^2} \int d^4x\, e^{i\kappa x} \langle B^b | \partial^\mu T\left(\bar{q}(x) P^a_\mu q(x), O^{MN}\right) -$$

$$- \delta(x^0) \left[\bar{q}(x) P^a_0 q(x), O^{MN}\right] | B^c \rangle, \tag{15}$$

or, in the soft-pion approximation ($\kappa \to 0$)

$$A^{MN} = \frac{i\sqrt{2}}{f_\pi} \int d^3x\, \langle B^b | \left[\bar{q}(x) P^a_0 q(x), O^{MN}\right]_{x^0=0} | B^c \rangle. \tag{16}$$

The equal-time commutator in (16) is easily calculated through the canonical rules:

$$C_6(x) = \left[\bar{q}(x) P^a_0 q(x), O^{MN}\right]_{x^0=0} =$$

$$= \delta^3(x) : q^+ \left[\gamma^0 P^a_0, \gamma^0 M\right] q\, \bar{q} N q : + \{M \leftrightarrow N\}. \tag{17}$$

Then allowing for the explicit form of P^a_μ, we find

$$A^{MN} = -\frac{i\sqrt{2}}{f_\pi} \langle B^b | : \bar{q} \left\{\gamma_5 \frac{\tau^a}{2}, M\right\} q\, \bar{q} N q : | B^c \rangle + \{M \leftrightarrow N\}. \tag{18}$$

Thus, the amplitude A^{MN} reduces to matrix elements of the local four-quark operators between one-particle hadron states.

In the course of calculation of such matrix elements it should be taken into account that initial and final hadron states always include the nonperturbative quark vacuum condensate, besides the quark determining the hadron states themselves (in our picture of the weak hadron interactions these are valence quarks). The existence of the condensate is exhibited immediately in the spontaneous breakdown of the chiral symmetry, and PCAC provides the known relation between the scalar density of this condensate and parameters of the explicit breaking of chiral symmetry (i.e., Lagrangian quark masses m_q):

$$2(m_u+m_d)\langle 0|\bar{q}q|0\rangle = -f_\pi^2 m_\pi^2 \qquad (\bar{q}q=\bar{u}u, \bar{d}d) \; . \tag{19}$$

Let us introduce the matrices Q and R through the Fierz transformation

$$M_{AB}\times N_{CD} = Q_{AD}\times R_{CB} \; . \tag{20}$$

Then, taking into account that

$$\begin{aligned}(P,M)_{AB}\times N_{CD} + M_{AB}\times(P,N)_{CD} &= \\ = (P,Q)_{AD}\times R_{CB} + Q_{AD}\times(P,R)_{CB} \; ,\end{aligned} \tag{21}$$

where P is a matrix and the brackets denote the commutator or anticommutator, the vacuum contributions may be singled out in the matrix elements of Eq. (18)

$$\begin{aligned}(A^{MN})^{vac} = &-\frac{i\sqrt{2}}{f_\pi}\Big(\langle 0|\bar{q}\{\gamma_5\tfrac{\tau^a}{2}, M\}q|0\rangle\langle B^b|\bar{q}Nq|B^c\rangle - \\ &-\langle 0|\bar{q}\{\gamma_5\tfrac{\tau^a}{2}, Q\}q|0\rangle\langle B^b|\bar{q}Rq|B^c\rangle\Big) + \{M\leftrightarrow N, Q\leftrightarrow R\}\end{aligned} \tag{22}$$

(see also ref. /18/).

Let us note, that the terms $\propto \langle 0|\bar{q}Kq|0\rangle\langle B^b|\bar{q}\{\gamma_5\frac{\tau^a}{2}, L\}|B^c\rangle$ ($K, L = M, N, Q, R$) do not contribute to A^{MN} because they correspond to the appearance in $\mathcal{H}^{PNC}$ of operators with the nonperturbative coefficient functions $\propto \langle 0|\bar{q}Kq|0\rangle$: such operators have the canonical dimension equal to 3 and are absorbed by the quark mass matrix counterms in the initial Lagrangian (see /12/).

Now the total amplitude A^{MN} is written down in the form

$$A^{MN} = (A^{MN})^{vac} - \frac{i\sqrt{2}}{f_\pi}\langle B^b|:\bar{q}\{\gamma_5\tfrac{\tau^a}{2}, M\}\bar{q}Nq:|B^c\rangle^q + \{M\leftrightarrow N\} \; , \tag{23}$$

where index q in the second matrix element denotes that the quark operators in it act only on quarks determining the proper baryon states.

To clarify to what mechanisms of the $\pi B'B$ interactions different contributions to A^{MN} correspond, we transform Eq. (23) into another form.

Vacuum contributions. Taking into account the relation

$$-\delta^3(x)\,\bar{q}\{\gamma_5\tfrac{\tau^a}{2}, K\}q = [\bar{q}(x)P_0^a q(x), \bar{q}Kq]_{x^0=0} \tag{24}$$

and using PCAC and the reduction formula for π "from right to left", we get

$$\begin{aligned}(A^{MN})^{vac} = \Big(&\langle \pi^a|\bar{q}Mq|0\rangle\langle B^b|\bar{q}Nq|B^c\rangle - \\ &-\langle \pi^a|\bar{q}Qq|0\rangle\langle B^b|\bar{q}Rq|B^c\rangle\Big)_{K\to 0} + \{M\leftrightarrow N, Q\leftrightarrow R\}.\end{aligned} \tag{25}$$

Hence, we see that A^{vac} corresponds to the known factorizable (F) contributions to A /26.17.18.28/ (Fig. 1a).

The "quark" contributions. Using the fact that in our approximation the operators O^{MN} coincide with the free ones (see sect. 2) we apply to the T - product in Eq. (15), the Wick theorem. With the equations of motion for the quark fields, we find

$$\partial^\mu T(:\bar{q}(x)P_\mu^a q(x), O^{MN}) = T_9(x,0) + T_6(x,0) + T_6'(x,0) + T_3(x,0), \tag{26}$$

$$T_9(x,0) = i:\bar{q}(x)P^a q(x)\,\bar{q}Mq\,\bar{q}Nq: , \tag{a}$$

$$T_6(x,0) = :[\bar{q}MS(-x)P^a q(x) + \bar{q}(x)P^a S(x)Mq]\bar{q}Nq: + \{M\leftrightarrow N\}, \tag{b}$$

$$T_6'(x,0) = -\delta^4(x):[\bar{q}M\gamma^0 P_0^a q(x) + \bar{q}(x)\gamma^0 P_0^a Mq]\bar{q}Nq: + \{M\leftrightarrow N\}, \tag{c}$$

$$\begin{aligned}T_3(x,0) = \partial^\mu[\mathrm{Tr}(S(-x)P_\mu^a S(x)M):\bar{q}Nq: - :\bar{q}MS(-x)P_\mu^a S(x)Nq:] &+ \\ + \{M\leftrightarrow N\} ,&\end{aligned} \tag{d}$$

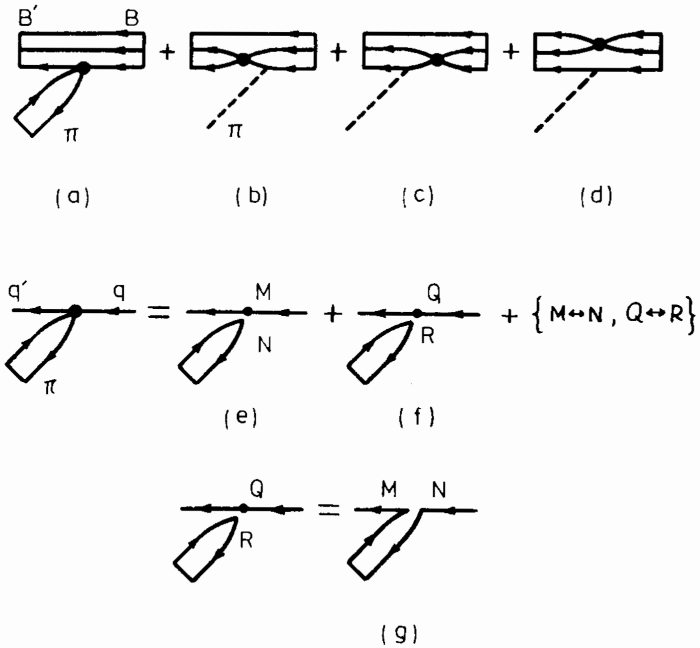

Fig. 1. Structure of the $\pi B'B$ vertex. The black circle denotes the effective Hamiltonian $\mathcal{H}^{PNC}$. (a) is factorizable (F) and (b)-(d) are nonfactorizable (NF) contributions; (e)-(g) is the interpetation of the F part of the vertex.

where $P^a = \gamma_5\{m_q, \frac{\tau^a}{2}\}$, $S(x) = (2\pi)^{-4}\int d^4p\,\exp(-ipx)\,(m_q - \not{p} - i\varepsilon)^{-1}$.

The lower indices of operators T in (26) denote their operator dimensions.

As $k \to 0$, the contributions T_i to A^{MN} for different i tend to zero independently of each other, in particular

$$\int d^4x\, e^{ikx}\langle B^b|\, T_6(x,0) + T_6'(x,0)\,|B^c\rangle_{k\to 0} \to 0\,. \tag{27}$$

Since from (17) and (26c) there follows

$$\int d^4x\, e^{ikx}\langle B^b|\, T_6'(x,0) - \delta(x^0)\, C_6(x)\,|B^c\rangle = 0\,, \tag{28}$$

allowing for Eq. (27) we get

$$(A^{MN})^q = -\frac{i\sqrt{2}}{f_\pi}\int d^4x\, e^{ikx}\langle B^b|\,T_6(x,0)\,|B^c\rangle^q_{k\to 0}\,. \tag{29}$$

Hence it follows that A^q corresponds to nonfactorizable (NF) contributions to A determined by the operator T_6 (Fig. 1 b,c) *). The NF contribution Fig. 1d is determined by the operator T_9 and vanishes as $k \to 0$; the operator T_3 gives (perturbative) contributions to A^F and vanishes in that limit, too.

3.2. Calculation of h_π

Now we shall apply the general formulas obtained in sect. 3.1 to calculate the constant h_π.

Factorizable (vacuum) contributions. From Eqs. (25) and (11), with taking into account the Fierz identities

$$(\gamma_\mu\gamma_5)_{\alpha\beta}\times(\gamma^\mu)_{\gamma\delta} \pm (\gamma_\mu)_{\alpha\beta}\times(\gamma^\mu\gamma_5)_{\gamma\delta} = -\left[\begin{pmatrix}\gamma_\mu\gamma_5\\ 2\gamma_5\end{pmatrix}_{\alpha\delta}\times\begin{pmatrix}\gamma^\mu\\ \delta\end{pmatrix}_{\gamma\beta} \pm \begin{pmatrix}\gamma_\mu\\ 2\delta\end{pmatrix}_{\alpha\delta}\times\begin{pmatrix}\gamma^\mu\gamma_5\\ \gamma_5\end{pmatrix}_{\gamma\beta}\right] \tag{30}$$

and the relation $\langle\pi|\bar q_i M q_j|0\rangle = \frac{1}{3}\delta_{ij}\langle\pi|\bar q_k M q_k|0\rangle$ (i, j, k are colour indices, a sum over k is understood) we find

$$\langle\pi^- p|\mathcal{H}^{PNC}|n\rangle^F = -2\sqrt{2}\,G\,(c_1^5 + \tfrac{1}{3}c_1^6)\,\langle\pi^-|\bar d\gamma_5 u|0\rangle\langle p|\bar u d|n\rangle\,. \tag{31}$$

To calculate the matrix elements of the quark densities in (31), we use the equations of motion for the quarks, then (see refs. /26,17,18,28/:

$$\langle\pi^-|\bar d\gamma_5 u|0\rangle = \frac{i}{\sqrt{(2\pi)^3\, 2k^0}}\,\frac{f_\pi m_\pi^2}{m_u + m_d}\,, \tag{32}$$

$$\langle p|\bar u d|n\rangle = \frac{M_\Xi - 3M_\Lambda + 2M_P}{m_u - m_s}\,\bar u_p u_n\,, \tag{33}$$

*) This conclusion confirms the arguments of the papers /7,30/.

where $m_q \equiv m_q(\mu_0)$ are the quark Lagrangian masses renormalized at the point $\mu = \mu_0$. To avoid small differences ($M_p - M_n$ and $m_u - m_d$) in (33), we use the SU(3) relation

$$\langle p | O_1 | n \rangle = -\sqrt{\tfrac{2}{3}} \left(2 \langle p | O_{1/2} | \Lambda \rangle - \langle \Lambda | O_{1/2} | \Xi^- \rangle \right) \qquad (34)$$

valid for any operators $O_{\Delta I = 1}$ and $O_{\Delta I = 1/2}$ which belong to the same octet.

From Eqs. (31) - (33) we obtain for h_π^F

$$h_\pi^F = -\frac{2\sqrt{2}}{3} \, G m_\pi^2 \left(3c_1^5 + c_1^6 \right) \frac{f_\pi}{m_u + m_d} \, \frac{M_\Xi - 3M_\Lambda + 2M_p}{m_u - m_s} \, . \qquad (35)$$

Thus, the calculation of h_π^F is reduced to the finding of the quark masses m_u, m_d, m_s.

The renormalized mass m_q as a function of the renormalization point μ in the lowest log approximation has the form (see, e.g., ref./25/)

$$m_q(\mu) = \bar{m}_q \left(\ln \frac{\mu}{\Lambda_3} \right)^{-\frac{4}{9}} , \qquad (36)$$

where $\bar{m}_q$ is a constant. From QCD sum rules for the axial current divergences it is known/25/: $\bar{m}_u = (7.6 \pm 2.2)$ MeV, $\bar{m}_d = (13.3 \pm 3.9)$ MeV, $\bar{m}_s = (260 \pm 80)$ MeV. Hence, according to Eq.(36) we have ($\Lambda_3 \simeq 100$ MeV) $m_u = (8.9 \pm 2.6)$ MeV, $m_d = (15.7 \pm 4.6)$MeV, $m_s = (306 \pm 94)$ MeV. We choose mean values from these intervals:

$$m_u = 8.9 \text{ MeV} , \qquad m_d = 15.7 \text{ MeV}, \qquad m_s = 306 \text{ MeV} . \qquad (37)$$

This choice is supported by the following:
(i) With the use of PCAC and the average SU(3) - splitting of hadron masses in the octets, in ref./31/ there have been obtained famous expessions for m_q : $m_u = (4.2 \text{ MeV})/Z$, $m_d = (7.5 \text{ MeV})/Z$, $m_s = (150 \text{ MeV})/Z$, where the constant $Z = \langle n | \bar{u} u | n \rangle = \langle p | \bar{d} d | p \rangle$. This constant is easily evaluated in quark models. In the MIT bag model $Z \simeq 0.48$, that provides the values of m_q practically coinciding with those of Eq. (37). (ii) The values of m_q of Eq. (37) correspond to the density of the quark condensate (see Eq. (19)) $\langle 0 | \bar{q} q | 0 \rangle_{\mu = 1 \text{GeV}} \simeq (-225 \text{ MeV})^3$ that agrees with the results of QCD sum rules for mesons (see ref./25/). (iii) According to (33)

and (37) $\langle p | \bar{u} d | n \rangle \simeq 0.51$, that is very close to the value calculated immediately in the MIT bag model : $\langle p | \bar{u} d | n \rangle_{bag} \simeq 0.48$. The proximity of these results also provides an argument for the self-consistency of our calculation scheme partly using the MIT bag model *) (see the calculation of the NF part of h_π below).

As a result, from Eqs. (35) and (37), for the values of C_1^5 and C_1^6 from table 2 we obtain

$$h_\pi^F \simeq 1.6 \times 10^{-7} . \qquad (38)$$

<u>Nonfactorizable ("quark") contributions</u>. As it follows from (17), the operators $C_6[O^{MN}]$ have the same symmetry properties as O^{MN}. Then (remind that the matrix elements of such operators are determined in our framework by only valence quarks) owing to the antisymmetry of quark wave functions in baryons, we have $\langle B^b | C_6[O^{27}] | B^c \rangle^q = \langle B^b | C_6[O^S] | B^c \rangle^q = 0$ **) and $\langle B^b | C_6[O^S] | B^c \rangle^q = -\langle B^b | C_6[O^6] | B^c \rangle^q$; besides, as all the components of O_1^A include the s-quark operators (see Appendix), $\langle p | C_6[O_1^A] | n \rangle^q = 0$. With this relations, from Eqs. (23) and (17), we find

$$h_\pi^{NF} = \frac{i}{f_\pi} \left(c_1^5 - c_1^6 \right) \langle p | : \bar{u} \gamma_\mu d \left(\bar{u} \gamma^\mu u + \bar{d} \gamma^\mu d + \bar{s} \gamma^\mu s \right) : - \qquad (39)$$
$$\left(\gamma_\mu \times \gamma^\mu \to \gamma_\mu \gamma_5 \times \gamma^\mu \gamma_5 \right) | n \rangle^q .$$

On the basis of Eq. (39) h_π^{NF} may be computed immediately in a quark model, e.g., in the MIT bag model, or connected by means of the SU(3) relation (34) with the amplitudes $A(\Lambda_-^0)^{NF}$ and $A(\Xi_-^-)^{NF}$. First consider the second, model-independent approach.

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*) Here it should be remembered, that the large ($\sim f_\pi / (m_u + m_d)$) value of the matrix element (32) is due to the pseudogoldstone nature of the pion. In the MIT bag model, where the chiral symmetry is violated explicity by the boundary conditions, $\langle \pi^- | \bar{d} \gamma_5 u | 0 \rangle_{bag} \sim f_\pi R_\pi m_\pi^2$,i.e.,it is by an order of magnitude as small as (32) (see ref./32/).

**) It is the so-called Pati-Woo argument/33/.

Since the operator $O^A_{1/2}$ gives no contribution to the r.h.s. of Eq. (34), we come to the following expression for h_π^{NF} :

$$h_\pi^{NF} = -\sqrt{\frac{2}{3}}\ \frac{C_1^5 - C_1^6}{C_{1/2}^5 - C_{1/2}^6}\ \left[2A(\Lambda^0_-)^{NF} - A(\Xi^-_-)^{NF} \right] . \tag{40}$$

This expression differs from analogous realizations of (34)(for review see refs./5,29/) in the following points. First, Eq. (40) is obtained for the <u>total</u> effective PNC Hamiltonian of the standard model ((10),(11)) *) without any *ad hoc* assumption on somewhat essential admixture of strange quarks in initial and final hadron states. Second, Eq. (40) connects only NF parts of the amplitudes. The impossibility of writing a sum rule of the type (40) for total (experimentally known) amplitudes $A(\Lambda^0_-)$ and $A(\Xi^-_-)$ is due to that combinations of the coefficients C^5 and C^6 in A^F and A^{NF} are different (cf. Eqs. (39) and (35)) and moreover that to $A(\Lambda^0_-, \Xi^-_-)^{NF}$ the contributions from the operators $O^{5,6}$ are complimented by those from O^{27}, O^S and O^A .

The following circumstance, however, hinders us to find the correct value of h_π^{NF} from Eq. (40): The r.h.s. of Eq. (40) is very sensitive to the values of the coefficients $C^5_{1/2}$ and $C^6_{1/2}$ (they enter both into the denominator of that relation and into $A^{NF} = A^{exp} - A^F$), and as we emphasized in sect. 2, errors of the leading log approximation for $C^{5,6}_{1/2}$ may essentially surpass those for other coefficients C_i^R . Therefore, to evaluate h_π^{NF} , we apply immediately to Eq. (39) and the MIT bag model. Then, we obtain

$$h_\pi^{NF} = 4\sqrt{2}\, G\ \frac{1}{f_\pi}\ (C_1^5 - C_1^6)\left(I_a - \frac{1}{3} I_b \right) , \tag{41}$$

*) Note, the relation (40) is a consequence of the choice of a low renormalization point μ and of the consideration of the different mass scales of m_c and $m_{u,d,s}$. Indeed, only in this case the $\Delta S = 1$ partners of $O_1^{5,6}$ appear in $\mathcal{H}^{PNC}_{\Delta S=1}$ (see sect.2) and the relations are valid, which are given prior to the Eq. (39).

where $I_a = \int_0^{R_N} r^2 dr \left[G^4(r) + F^4(r) \right]$, $I_b = 2\int_0^{R_N} r^2 dr\, G^2(r) F^2(r)$

$G(r) = J_0\left(\frac{\omega}{R} r\right)$ and $F(r) = -J_1\left(\frac{\omega}{R} r\right)$ are radial parts of a "large" and a "small" component of the quark bispinor in the MIT bag model ($m_u = m_d = 0$, $\omega \simeq 0.24$). At $R_N \simeq 1$ fm (the nucleon radius in the model) $I_a \simeq 1.43 \times 10^{-3}$ GeV^3, $I_b \simeq 0.65 \times 10^{-3}$ GeV^3 , and we get

$$h_\pi^{NF} \simeq -0.31 \times 10^{-7} . \tag{42}$$

<u>**The total value of h_π.**</u> Now we have found all the components of h_π . According to Eqs. (23), (38) and (42) we get

$$h_\pi = h_\pi^F + h_\pi^{NF} \simeq 1.3 \times 10^{-7} . \tag{43}$$

Separating the contributions to h_π from the charged and neutral currents (see Eq. (4)), we find

$$h_\pi = h_\pi(C.C.) + h_\pi(N.C.) \simeq (0.0 + 1.3)\times 10^{-7} \tag{44}$$

Hence it follows that h_π is almost completely determined by the neutral currents (see also ref./28/). We give also the value of h_π for the bare effective Hamiltonian $\mathcal{H}^{PNC}_{\alpha_S=0}$.
In this case $C_1^5 = 0$, $C_1^6 = -\frac{1}{6} S_W^2$ and

$$(h_\pi)_{\alpha_S=0} \simeq 0.48 \times 10^{-7} . \tag{45}$$

Thus, the quark-gluon interactions increase h_π about 3 times (while $h_\pi^F/(h_\pi^F)_{\alpha_S=0} \simeq 7.3$, $h_\pi^{NF}/(h_\pi^{NF})_{\alpha_S=0} \simeq -1.2$).
From Eqs. (38) and (42) it is seen that in h_π the F part dominates in agreement with the results of papers/17,18,28/. However, our value of h_π is at least 2 times as small as previous estimates of h_π within the standard model. This point is explained by the following: First, we have used larger values of m_q than have been used formerly, that is consistent with the modern data and with our application of the MIT bag model (see point (iii) after Eq. (37)). Second, for the calculation of h_π^{NF} we have consistently used the MIT bag model and do not apply to relations of the type (40), the incorrect use of which may lead to (erroneous) large positive values of h_π^{NF} . In the conclusion of this section, we provide the interval of values of h_π arising due to the uncertainty of our knowledge of m_q . For $m_u + m_d = (17 \to 32)$MeV (i.e.,

$\langle 0|\bar q q|0\rangle_{\mu=1\,GeV} = -(200 \leftarrow 250)^3\ \mathrm{MeV}^3)$ and $m_s - m_u = (205 \to 390)$ MeV,

we find *)

$$h_\pi = (0.62 \leftarrow 3.0) \times 10^{-7}. \qquad (46)$$

Thus, by the experimental determination of h_π the values of quark mass (or the quark condensate density) could be defined more accurately, while the observation of h_π on the level of $h_\pi^{b.v.} \simeq 4.6 \times 10^{-7}$ would give explicit evidence for some exotic contributions to the PNC NN forces (e.g. /9/) or for insufficiency of the leading log approximation for $\mathcal{H}^{PNC}$.

The experimental consequences of the values (43) - (45) will be descussed in a concluding paper(III).

We are grateful to M.A.Shifman for valuable discussions.

Appendix. Explicit form of the operators of the effective Hamiltonian $\mathcal{H}^{PNC}$ (10),(11).

Notation: $\bar q_1 q_2 \bar q_3 q_4 \equiv :\bar q_{1i}\gamma_\mu\gamma_5 q_{2i}\ \bar q_{3j}\gamma^\mu q_{4j}:$, a sum over color indeces i,j = 1,2,3 is understood.

$\Delta S = 1$

$$O^{27}_{3/2,1/2} = -2(\bar u u \bar d s + \bar d s \bar u u + \bar u s \bar d u + \bar d u \bar u s) + \bar d d \bar d s + \bar d s \bar d d + \bar d s \bar s s + \bar s s \bar d s + h.c.,$$

$$O^S_{1/2} = \bar u u \bar d s + \bar d s \bar u u + \bar u s \bar d u + \bar d u \bar u s + 2(\bar d d \bar d s + \bar d s \bar s d + \bar d s \bar d d + \bar s s \bar d s) + h.c.,$$

$$O^A_{1/2} = -\bar u u \bar d s - \bar d s \bar u u + \bar u s \bar d u + \bar d u \bar u s + h.c.,$$

*) It is to be remembered, however, that the deviation of the values of m_q from (37) violates the self-consistency of calculation of h_π with the use of the MIT bag model, so $\langle p|\bar u d|n\rangle = (0.76 \leftarrow 1.4)\langle p|\bar u d|n\rangle_{bag}$.

$$O^5_{1/2} = -(\bar u_i u_j + \bar d_i d_j + \bar s_i s_j)\bar d_j s_i + \bar d_i s_j(\bar u_j u_i + \bar d_j d_i + \bar s_j s_i) + h.c.,$$

$$O^6_{1/2} = -(\bar u u + \bar d d + \bar s s)\bar d s + \bar d s(\bar u u + \bar d d + \bar s s) + h.c. \qquad (A1)$$

$\Delta S = 0$

$$O^{27}_2 = -(\bar u u \bar u u + \bar d d \bar d d) + \bar u u \bar d d + \bar d d \bar u u + \bar u d \bar d u + \bar d u \bar u d,$$

$$O^{27}_1 = \bar u u \bar u u - \bar d d \bar d d - (\bar u u \bar s s + \bar s s \bar u u + \bar u s \bar s u + \bar s u \bar u s) + \bar d d \bar s s + \bar s s \bar d d + \bar d s \bar s d + \bar s d \bar d s,$$

$$O^{27}_0 = 2(\bar u u \bar u u + \bar d d \bar d d + 3\bar s s \bar s s) + \bar u u \bar d d + \bar d d \bar u u + \bar u d \bar d u + \bar d u \bar u d - 3(\bar u u \bar s s + \bar s s \bar u u + \bar d d \bar s s + \bar s s \bar d d + \bar u s \bar s u + \bar s u \bar u s + \bar d s \bar s d + \bar s d \bar d s),$$

$$O^S_1 = 2(\bar u u \bar u u - \bar d d \bar d d) + \frac{1}{2}(\bar u u \bar s s + \bar s s \bar u u + \bar u s \bar s u + \bar s u \bar u s) - \frac{1}{2}(\bar d d \bar s s + \bar s s \bar d d + \bar d s \bar s d + \bar s d \bar d s),$$

$$O^A_1 = -\frac{1}{2}(\bar u u \bar s s + \bar s s \bar u u + \bar d s \bar s d + \bar s d \bar d s) + \frac{1}{2}(\bar d d \bar s s + \bar s s \bar d d + \bar u s \bar s u + \bar s u \bar u s),$$

$$O^5_1 = \bar u_i u_j \bar d_j d_i - \bar d_i d_j \bar u_j u_i + \frac{1}{2}(\bar u_i u_j \bar s_j s_i - \bar s_i s_j \bar u_j u_i - \bar d_i d_j \bar s_j s_i + \bar s_i s_j \bar d_j d_i),$$

$$O^6_1 = \bar u u \bar d d - \bar d d \bar u u + \frac{1}{2}(\bar u u \bar s s - \bar s s \bar u u - \bar d d \bar s s + \bar s s \bar d d),$$

$$O^S_0 = \frac{1}{\sqrt{3}}\Big[2(\bar u u \bar u u + \bar d d \bar d d - 2\bar s s \bar s s) + \bar u u \bar d d + \bar d d \bar u u + \bar u d \bar d u + \bar d u \bar u d - \frac{1}{2}(\bar u u \bar s s + \bar s s \bar u u + \bar d d \bar s s + \bar s s \bar d d + \bar u s \bar s u + \bar s u \bar u s + \bar d s \bar s d + \bar s d \bar d s)\Big],$$

$$O^A_0 = \frac{1}{\sqrt{3}}\Big[-\bar u u \bar d d - \bar d d \bar u u + \bar u d \bar d u + \bar d u \bar u d + \frac{1}{2}(\bar u u \bar s s + \bar s s \bar u u + \bar d d \bar s s + \bar s s \bar d d) - \frac{1}{2}(\bar u s \bar s u + \bar s u \bar u s + \bar d s \bar s d + \bar s d \bar d s),$$

$$O_0^5 = \frac{\sqrt{3}}{2}\left(\bar{u}_i u_j \bar{s}_j s_i - \bar{s}_i s_j \bar{u}_j u_i + \bar{d}_i d_j \bar{s}_j s_i - \bar{s}_i s_j \bar{d}_j d_i\right),$$

$$O_0^6 = \frac{\sqrt{3}}{2}\left(\bar{u}u\bar{s}s - \bar{s}s\bar{u}u + \bar{d}d\bar{s}s - \bar{s}s\bar{d}d\right),$$

$$O^{1S} = 2\left(\bar{u}u\bar{u}u + \bar{d}d\bar{d}d + \bar{s}s\bar{s}s\right) + \bar{u}u\bar{d}d + \bar{d}d\bar{u}u + \bar{u}u\bar{s}s + \bar{s}s\bar{u}u +$$

$$+ \bar{d}d\bar{s}s + \bar{s}s\bar{d}d + \bar{u}d\bar{d}u + \bar{d}u\bar{u}d + \bar{u}s\bar{s}u + \bar{s}u\bar{u}s + \bar{d}s\bar{s}d + \bar{s}d\bar{d}s,$$

$$O^{1A} = -\left(\bar{u}u\bar{d}d + \bar{d}d\bar{u}u + \bar{u}u\bar{s}s + \bar{s}s\bar{u}u + \bar{d}d\bar{s}s + \bar{s}s\bar{d}d\right) +$$

$$+ \bar{u}d\bar{d}u + \bar{d}u\bar{u}d + \bar{u}s\bar{s}u + \bar{s}u\bar{u}s + \bar{d}s\bar{s}d + \bar{s}d\bar{d}s. \quad \text{(A2)}$$

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Дубовик В.М., Зенкин С.В. E2-83-611

Самосогласованный расчет слабых констант в несохраняющих четность ядерных силах.
Эффективный НЧ гамильтониан в $SU(2)_L xU(1)xSU(3)_c$. НЧ в πNN вершине.

На основе полного эффективного гамильтониана несохраняющих четность /НЧ/ адрон-адронных взаимодействий, найденного в стандартной модели $SU(2)_L x xU(1)xSU(3)_c$ во всех порядках главных логарифмов, с учетом различия масштабов кварковых масс ($m_c >> m_{u,d,s}$) рассматривается НЧ πNN вершина, генерирующая дальнодействующую часть НЧ ядерных сил. Мы анализируем природу и методы расчета различных вкладов в эту вершину, обращая особое внимание на возможные артефакты этих методов. В рамках самосогласованной расчетной схемы, частично включающей модель массачусетского мешка, рассчитывается полное значение константы h_π, определяющей эту вершину. Наше значение h_π ($\approx 1.3x10^{-7}$) в 2÷4 раза меньше предыдущих оценок и не противоречит экспериментальным данным.

Работа выполнена в Лаборатории теоретической физики ОИЯИ.

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Dubovik V.M., Zenkin S.V. E2-83-611

Self-Consistent Calculation of the Weak Constants in the Parity Nonconserving Nuclear Forces.
Effective PNC Hamiltonian in $SU(2)_L xU(1)xSU(3)_c$.
PNC in the πNN Vertex

On the basis of the total effective Hamiltonian of the parity nonconserving /PNC/ hadron-hadron interactions found within the standard model $SU(2)_L x xU(1)xSU(3)_c$ in all orders of the leading logarithms allowing for the difference of quark mass scales ($m_c >> m_{u,d,s}$) we consider the PNC πNN vertex generating the long-range part of the PNC nuclear forces. We analyse the origin and the methods of calculation of various contributions to this vertex with a special attention to possible artifacts of these methods. Within the self-consistene calculational framework partly including the MIT bag model we evaluate the total value of the constant h_π determining the PNC πNN vertex. Our value of h_π ($\approx 1.3x10^{-7}$) is 2÷4 times as small as previous estimates and does not contradict the experimental data.

The investigation has been performed at the Laboratory of Theoretical Physics, JINR.

Communication of the Joint Institute for Nuclear Research. Dubna 1983