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

PNC in the ρNN and ωNN Vertices

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1. INTRODUCTION

This paper is a sequel to our work [1] (further referred to as I) on a systematic consideration in the standard model $SU(2)_L \times U(1) \times SU(3)_c$ of parity nonconserving (PNC) MNN vertices $(M = \pi, \rho, \omega)$ which determine the PNC NN potential. We calculate here constants $h_\rho^{0,1,2}$ and $h_\omega^{0,1}$ (see I, sect. 1). The invariant amplitude of the PNC transition $N \to N'V$ $(V = \rho, \omega)$ with the change of isospin by $i = 0,1,2$ has the form

$$< VN'|\mathcal{H}_i^{PNC}|N> = \frac{\epsilon^{*\mu}}{\sqrt{(2\pi)^3\, 2k^0}} M_\mu^i(k) , \tag{1}$$

$$M_\mu^i(k) = \bar{N}'[h^i_{VN'N}(k^2)\gamma_\mu + if^i_{VN'N}(k^2)k_\mu + g^i_{VN'N}(k^2)\sigma_{\mu\nu}k^\nu]\gamma_5 N ,$$

where ϵ^μ is the polarization vector of V. The constants we are interested in are related to the form factors $h(k^2)$ as follows

$$h_\rho^{0,1} = h_{\rho pp}^{0,1}(0), \quad h_\rho^2 = \sqrt{6}\, h_{\rho pp}^2(0), \quad h_\omega^{0,1} = h_{\omega pp}^{0,1}(0) . \tag{2}$$

First estimates of the constants h_V [2] were based on the phenomenological VA form of the effective Hamiltonian and on the assumption of factorization of the amplitude (1):

$$< VN'|\mathcal{H}^{PNC}|N> \simeq \sqrt{2}\, G \sum_{a,b} C_{ab} < V|V_\mu^a|0> < N'|A^{b\mu}|N> . \tag{3}$$

In ref. [3] this approximation was substantiated in the framework of current algebra. In refs. [4,5] it has been observed that within quark models the right-hand side of the expression (3) is not invariant under the Fierz transformation and determines not all contributions to the factorizable (F) part of the amplitude (1). This shortcoming of the approximation (3) was eliminated in a modified factorization approach [5,6] based on the field theoretical consideration of the PNC VNN vertices. However, as the analysis performed in a number of papers (see, e.g., refs. [6,7]) has shown, even with gluon corrections the constants h_V^F do not provide the agreement of theoretical and experimental results.

Another approach to the calculation of the constants h_V in the Cabibbo weak interaction model has been proposed in ref.[8]. There the amplitudes (1) were expanded over irreducible representations of the group $SU(6)_W$ and were related to S-wave amplitudes of nonleptonic decays of hyperons. Values obtained within this approach for h_V turn out to be very different from values of h_V^F calculated in the Cabibbo model. This situation was explained in ref.[9], in which $SU(6)_W$ results of ref.[8] were interpreted in terms of quark diagrams: values of h_V obtained in[8] correspond to a nonfactorizable (NF) parts of matrix elements $\langle VN'|\mathcal{H}^{PNC}|N\rangle$ and should be summed with h_V^F. To calculate the NF contributions to the constants h_V in the standard model, in[9] the $SU(6)_W$ was completed by a nonrelativistic quark technique. In this approach, the NF parts of the constants were expressed in terms of three parameters ($\tilde{b}_t$, $\tilde{b}_V$, $\tilde{c}_V$) which were determined from the known S-wave amplitudes of nonleptonic decays (the parameter $\tilde{c}_V$ was interpreted as a contribution of the quark sea to matrix elements). Besides, to get agreement with experimental data, further factors were introduced to h_V^{NF}. Values of h_M obtained in[9] are known as the "best values" ($h_M^{b.v.}$).

A direct calculation of h_ρ^{NF} was attempted in refs.[10,11]. In these papers the NF contributions to h_ρ were approximated by pole contributions of P-odd nucleon resonances $N^*(1/2^-)$, and matrix elements of PNC transitions $N^* \to N$ were calculated in the MIT bag model. Values of h_ρ^{NF} found in[11] have the same signs as $(h_\rho^{NF})^{b.v.}$, but in absolute value they are ~ 1.5 times as small as the latter. As a result, the constant h_ρ^0 in[11] practically disappears because the F and NF contributions cancel out. Results of[11], however, should not be considered completive: the MIT bag model may be inadequate for the calculation of nonstatic matrix elements $\langle N|\mathcal{H}^{PNC}|N^*\rangle$.

In this paper we shall find the NF contributions to constants h_V in the scheme in which the matrix elements of the operator part of the effective Hamiltonian are defined by only valence quarks (see I, sects. 2,3) using the approximate SU(6) symmetry of the matrix elements $\langle MB'|\mathcal{H}^{PNC}|B\rangle^{NF}$ (B, B' are baryons). h_V^{NF} will be connected with the know NF contributions to the S-wave amplitudes $B \to B'\pi$ and calculated in the MIT bag model. This approach gives for h_V values close to $h_V^{b.v.}$, and at our value $h_\pi \simeq \frac{1}{3} h_\pi^{b.v.}$ (see I, sect. 3) allows us to come up to the experimental results on PNC low-energy NN interactions without arbitrary (fitting) parameters.

In sect.2, we shall consider the overall structure of PNC VNN vertices and calculate the F parts of the constants h_V. In sect.3 we shall present the calculation of NF contribution to h_V, and in sect. 4 the results will be discussed.

2. OVERALL STRUCTURE OF PNC VNN VERTICES. THE F PARTS OF CONSTANTS h_V

Like in I, sect. 3, we shall write the effective Hamiltonian $\mathcal{H}^{PNC}$ in the form

$$\mathcal{H}^{PNC} = \sqrt{2}\, G \sum_{M,N} C^{MN} \mathcal{O}^{MN}, \tag{4}$$

$$\mathcal{O}^{MN} = :\bar{q} M q \bar{q} N q: , \tag{5}$$

and consider the partial amplitude

$$\langle VN'|\mathcal{O}^{MN}|N\rangle = \frac{\epsilon^{*\mu}}{\sqrt{(2\pi)^3\, 2k^\circ}} M_\mu^{MN}. \tag{6}$$

According to the accepted picture of PNC hadron-hadron interactions (see I) the matrix elements (6) are determined by the valence quarks, whereas contributions to the total amplitude

$$M_\mu = \sqrt{2}\, G \sum_{M,N} C^{MN} M_\mu^{MN} \tag{7}$$

of the nonvalence quarks are taken into account in the coefficient functions of the effective Hamiltonian - C^{MN}.

Let us apply to V in (6) the reduction formula. Using then the standard representation for the interpolating field of ρ- and ω-mesons

$$V_\mu^a = \frac{f_\rho}{m_\rho^2} \bar{q}_i \gamma_\mu \frac{\tau^a}{2} q_i \tag{8}$$

($a = 0,1,2,3$, $\vec{V}_\mu = \vec{\rho}_\mu$, $V_\mu^0 = \omega_\mu$, $\tau^0 = 1$, $f_\rho \simeq 5.1$ is the $\rho^0 \to e^+e^-$ decay constant; a sum over the colour index $i = 1,2,3$ is carried out) and the Wick theorem, we come to the following expressions [5,6]

$$\langle VN'|\mathcal{O}^{MN}|N\rangle = \langle VN'|\mathcal{O}^{MN}|N\rangle^F + \langle VN'|\mathcal{O}^{MN}|N\rangle^{NF}, \tag{9}$$

$$\begin{aligned} \langle VN'|\mathcal{O}^{MN}|N\rangle^F = {} & \langle V|\bar{q}Mq|0\rangle \langle N'|\bar{q}Nq|N\rangle - \\ & - \langle V|\bar{q}Qq|0\rangle \langle N'|\bar{q}Rq|N\rangle + \{M \leftrightarrow N,\ Q \leftrightarrow R\}, \end{aligned} \tag{10}$$

$$\langle VN'|\mathcal{O}^{MN}|N\rangle^{NF} = \frac{\epsilon^{*\mu}}{\sqrt{(2\pi)^3\, 2k^\circ}} \frac{f_\rho}{m_\rho^2} (-k^2 + m_V^2)\, i \int d^4x\, e^{ikx} \times \tag{11}$$

$$\times \langle N'|T_{9\mu}(x,0) + T_{6\mu}(x,0)|N\rangle .$$

Here

$$Q_{AB} \times R_{CD} = M_{AD} \times N_{CB} ,$$

$$T_{9\mu}(x,0) = :\bar{q}(x)\gamma_\mu \frac{\tau^a}{2}\bar{q}(x)\, qMq\bar{q}Nq: , \qquad (12)$$

$$T_{6\mu}(x,0) = -i: [\,\bar{q}MS(-x)\gamma_\mu \frac{\tau^a}{2} q(x) + \bar{q}(x)\gamma_\mu \frac{\tau^a}{2} S(x)\,Mq\,]\bar{q}Nq: + \{M \leftrightarrow N\}; \qquad (13)$$

for other notation see I, sec. 3.

We see that the amplitude $\langle VN'|\mathcal{C}^{MN}|N\rangle$ is represented by the sum of the F and NF parts invariant with respect to the Fierz transformations and defined by the expressions (10) and (11) and has the same quark structure as the amplitude $\langle \pi N'|\mathcal{C}^{MN}|N\rangle$ (I, sect. 3). Diagrams for the expressions (10) and (11) are completely analogous to the diagrams for the PNC πNN vertex (Fig.1 in I). From comparison of (9)-(11) with (3) it is also seen that the approximation of factorization (3), and consequently approaches based on current algebra take into account only "separable" parts of matrix elements $\langle VN'|\mathcal{C}^{MN}|N\rangle^F$ corresponding to the diagram (e) of Fig.1 in I (with the change $\pi \to V$).

The expression (10) allows us to find the F parts of constants h_V without any model constructions. Indeed, for valence quark nonzero matrix elements in the r.h.s. of (10) are reduced straighforward to the experimentally known matrix elements of hadron currents

$$\langle V^a|\bar{q}_i\gamma_\mu \frac{\tau^b}{2} q_j|0\rangle = \frac{1}{3}\delta_{ij}\,\delta^{ab}\frac{\epsilon^{*\mu}}{\sqrt{(2\pi)^3\, 2k^o}}\frac{m_\rho^2}{f_\rho} , \qquad (14)$$

$$\langle N'|\bar{q}\gamma_\mu\gamma_5 \frac{\vec{\tau}}{2} q|N\rangle = g_A \bar{N}'\gamma_\mu\gamma_5\frac{\vec{\tau}}{2}N , \qquad (15)$$

$$\langle N'|\bar{q}\gamma_\mu\gamma_5 \frac{\tau^o}{2} q|N\rangle = \zeta g_A \bar{N}'\gamma_\mu\gamma_5\frac{\tau^o}{2}N . \qquad (16)$$

In the expressions (15), (16) $g_A \simeq 1.25$ is the axial constant of the neutron β-decay, $\zeta = (-D + 3F)/g_A$; from the experimental data on lepton decays of $1/2^+$-baryons $D \simeq 0.80$, $F \simeq 0.45$ /12/ , hence $\zeta \simeq 0.44$.

Using the explicit form of the operators of $\mathcal{H}^{PNC}$ (see Appendix in I), the Fierz identity (I.30) and formulae (10), (14)-(16), we get

$$(h_\rho^o)^F = \frac{2}{3}G_V(2c_o^{27} + \frac{2}{\sqrt{3}}c_o^S + 2c^{1S} + \frac{1}{\sqrt{3}}c_o^A + c^{1A}) , \qquad (17)$$

$$(h_\rho^1)^F = \frac{1}{3}G_V\zeta(4c_1^{27} + 8c_1^S - c_1^5 - 3c_1^6) , \qquad (18)$$

$$(h_\rho^2)^F = -\sqrt{6}\,\frac{8}{3}G_V c_2^{27} , \qquad (19)$$

$$(h_\omega^o)^F = \frac{2}{3}G_V\zeta(6c_o^{27} + 2\sqrt{3}\,c_o^S + 6c^{1S} - \frac{1}{\sqrt{3}}c_o^A - c^{1A}) , \qquad (20)$$

$$(h_\omega^1)^F = \frac{1}{3}G_V(4c_1^{27} + 8c_1^S + c_1^5 + 3c_1^6) , \qquad (21)$$

where

$$G_V = \sqrt{2}\,G m_\rho^2 \frac{g_A}{f_\rho} \simeq 2.5\times 10^{-6} .$$

3. CALCULATION OF h_V^{NF}

As can be seen from the derivation of the expressions (11)-(13), the amplitudes $\langle VN'|\mathcal{H}^{PNC}|N\rangle^{NF}$ are not reduced to one-particle matrix elements of local operators, and therefore, the NF contributions to constants h_V cannot be calculated directly, for instance, with the use of the MIT bag model. In I, sect.3 we have shown that to such matrix elements the amplitudes $\langle \pi B'|\mathcal{H}^{PNC}|B\rangle^{NF}$ are reduced. Therefore we shall achieve our aim if we shall determine the functional structure of both the amplitudes in terms of the same parameters*. Indeed, then the parameters will be fixed by the structure of matrix elements $\langle \pi B'|\mathcal{H}^{PNC}|B\rangle^{NF}$ known in the MIT bag model, and therefore the NF parts of all the constants h_M will be determined through the overlapping integral of the bag quark wave functions.

The functional structure of the amplitudes $A_{MB'B}^{NF} = \sqrt{(2\pi)^3\, 2k^o}\,\langle MB'|\mathcal{H}^{PNC}|B\rangle^{NF}$ may be obtained by using the recipe

*By the functional structure of an amplitude we mean its representation by a linear combination of the coefficient functions c_i^R.

proposed in [9] (see also ref. [13]). In the picture we have accepted for hadron-hadron interactions it is reduced to the following two approximations:

1. The representation of the amplitudes $A^{NF}_{MB'B}$ via the quark transition amplitudes

$$A^{NF}_{MB'B} = b_{MB'B}\left[\langle (q\bar{q})_M (qqq)_{B'} | \mathcal{H}^{PNC} | (qqq)_B \rangle^{NF} + \right. \qquad (22)$$

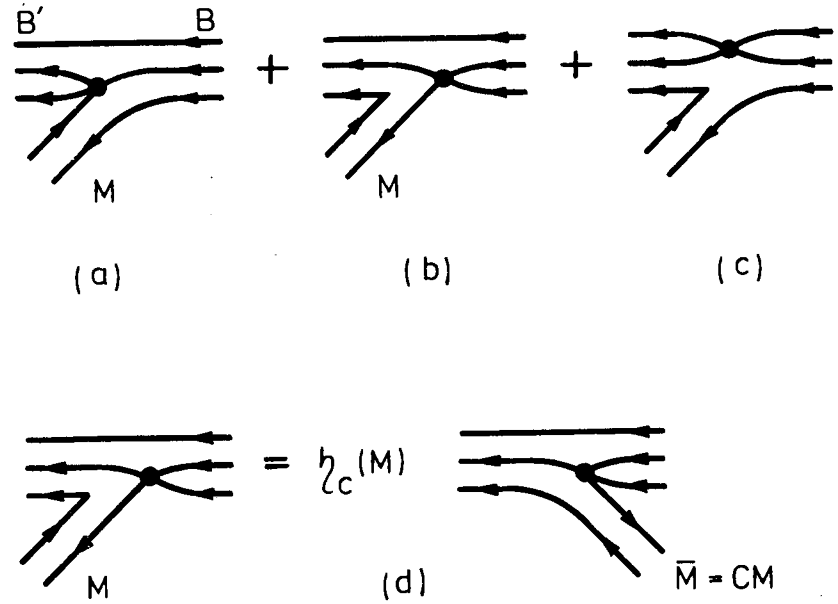
$$\left. + \eta_C(M) \langle (qqq)_{B'} | \mathcal{H}^{PNC} | (qqq)_B \, C (q\bar{q})_N \rangle^{NF} \right],$$

where $b_{MB'B}$ are parameters; $|(q\bar{q})_M\rangle$ and $|(qqq)_B\rangle$ are normalized states with quantum numbers of the meson M and baryon B constructed from the quark and antiquark and three quarks, respectively, with momenta $\vec{p}_q = \vec{p}_{\bar{q}} = 0$; $\eta_C(M)$ is the charge phase of the isotopic multiplet which includes M$[\eta_C(\pi) = 1, \eta_C(\rho) = = \eta_C(\omega) = -1]$, **C** is the charge-conjugation operator.

2. The SU(6) symmetry of the amplitudes $A^{NF}_{MB'B}$.

Let us comment in brief points 1 and 2. According to 1 the spin, flavour and colour structure of the amplitudes are determined by the quark matrix elements of the effective Hamiltonian $\mathcal{H}^{PNC}$. Diagrams of the quark transitions are shown in the Figure. Diagrams a), b) and c) correspond to the NF amplitudes b), c) and d) of Fig.1 in I. The contribution of diagram a) to $A^{NF}_{MB'B}$ is described by the first term in (22). The transitions b) and c) include the creation of a quark-antiquark pair correlated with the PNC 4q-interaction, and that is why they cannot be calculated straightforward. Therefore the contribution of the diagram b) into $A^{NF}_{MB'B}$ is replaced, owing to crossing symmetry, by the contribution from the $\bar{M}$-meson absorption diagram d) that does not contain the $q\bar{q}$ vertex. This contribution corresponds to the second term in (22). As to the diagram c), according to the relations (I.15), (I.26) and (11)-(13) it may be neglected at all. Indeed, this diagram corresponds to the matrix elements of operators T_9. The amplitude $A_{\pi B'B}(T_9)$, as is shown in I, sect.3, disappears as $k \to 0$, and the amplitude $A_{VN'N}(T_9)$ can be neglected because of the vector current conservation that may be verified by considering the longitudinal part of that amplitude. It then follows that the dominant contribution to $A^{NF}_{MB'B}$ comes from diagrams a) and b).

Parameters $b_{MB'B}$ in (22) mean the corresponding to diagram a) overlapping integrals of spatial parts of the hadron wave functions $\psi_M(\vec{x}_1 - \vec{x}_2)$, $\Psi_{B'}(\vec{y}_1 - \frac{\vec{y}_2 + \vec{y}_3}{2}, \vec{y}_2 - \vec{y}_3)$ and $\Psi_B(\vec{z}_1 - \frac{\vec{z}_2 + \vec{z}_3}{2}, \vec{z}_2 - \vec{z}_3)$. Owing to approximation 2 these inte-



The PNC transitions $(qqq)_B \to (q\bar{q})_M + (qqq)_{B'}$. The black circle denotes the effective Hamiltonian $\mathcal{H}^{PNC}$; (d) is the diagram corresponding to the absorption of the meson $\bar{M}$ = **C**M.

grals for all mesons M from the 35-plet and all baryons B,B' from the 56-plet have the same value, i.e., $b_{MB'B} = b$. Note that the "integral" approximation 2 is weaker than the approximation of SU(6) symmetry of the hadron wave functions *.

Thus, in the approximations 1 and 2 the amplitudes $A^{NF}_{MB'B}$ are determined by the quark amplitudes a), d) and one parameter b .

To control the accuracy with which the expression (22) determines the functional structure of $A^{NF}_{\pi B'B}$, and consequently, the parameter b , we shall consider as a benchmark the three amplitudes $A^{NF}_{\pi^- pn}$, $A^{NF}_{\pi^- p\Lambda}$ and $A^{NF}_{\pi^- \Lambda\Xi^-}$ (they enter into the sum rule (I.40)).

The constants h^{NF}_M and $A(B^{0,-}_-)^{NF}$ we are interested in are defined by the matrix elements (see Eqs. (1), (2) and (I.12)):

*So, the well-known violation of SU(6) symmetry of the meson wave functions $|\psi_M(0)|^2 / |\psi_{M'}(0)|^2 \sim m_M / m_{M'}$ (see [13]) in 2 is smoothed by integration of the functions $\psi_M(\vec{x}_1 - \vec{x}_2)$ with functions $\Psi_B \cdot \Psi_{B'}$.

$$(h_V^i)^{NF} = \sqrt{(2\pi)^3 2k^\circ} < V^\circ_{(o)} p_\uparrow | \mathcal{H}_i^{PNC} | p_\uparrow >^{NF}, * \quad (23)$$

$$A(B^{o,-}_{-})^{NF} = \frac{1}{i}\sqrt{(2\pi)^3 2k^\circ} < \pi^- B'_\uparrow | \mathcal{H}^{PNC} | B^{o,-}_\uparrow >^{NF}, \quad (24)$$

where $V^\circ_{(o)}$ is either ρ°- or ω- meson with the spin projection $s_3 = 0$, $B_\uparrow$ is the baryon with the spin projection $s_3 = 1/2$; in this notation $h_\pi = A(n^\circ_-)$. The results of calculations (see Appendix)** are as follows

$$(h^\circ_\rho)^{NF} = -16G(\frac{1}{\sqrt{3}} c_0^A + c^{1A})\, b/\sqrt{6}, \quad (25)$$

$$(h^1_\rho)^{NF} = -\frac{8}{3} G(c_1^5 - c_1^6)\, b/\sqrt{6}, \quad (26)$$

$$(h^2_\rho)^{NF} = (h^{0,1}_\omega)^{NF} = 0, \quad (27)$$

$$(h_\pi)^{NF} = 4\sqrt{2}\, G(c_1^5 - c_1^6)\, b/\sqrt{6}, \quad (28)$$

$$A(\Lambda^\circ_-)^{NF} = -\sqrt{12}\, G[\, c^A_{1/2}\, 2b/\sqrt{6} + (c^5_{1/2} - c^6_{1/2})\, b/\sqrt{6}\,], \quad (29)$$

$$A(\Xi^-_-)^{NF} = -4\sqrt{12}\, G\, c^A_{1/2}\, b/\sqrt{6}. \quad (30)$$

As is seen from (25)-(30), the symmetric operators $\mathcal{O}^{27}$, $\mathcal{O}^S$ and $\mathcal{O}^{1S}$ (see I, sect.2) do not contribute to the NF parts of the amplitudes that is a consequence of the antisymmetry of the quark wave functions in the baryons (the so-called Pati-Woo argument[14]; see also Appendix). For this reason the constant h^2_ρ determined by the operator $\mathcal{O}^{27}_2$ has only the F part. The constants $(h^{0,1}_\omega)^{NF}$ vanish since the contributions of $u\bar{u}$ and $d\bar{d}$ components of the vector meson wave functions to the matrix elements (23) are equal in absolute value and opposite in sign. Note that expressions (28-30) satisfy the sum rule (I.40), and from (26) and (28) the relation

* For $i = 2$ $h^i_V/\sqrt{6}$.

**Note that the expressions for matrix elements $< MB' | \mathcal{O}^{MN} | B >^{NF}$ obtained in Appendix and in work[9] coincide if in the latter we set $\tilde{b}_t = -\tilde{b}_v = 4\sqrt{3}\, b$ and $\tilde{c}_v = 0$. (Besides, in Table III of[9] one should change the sigh of the NF part of the matrix element $< \pi^- p | \mathcal{O}_5 | n >$).

$$(h^1_\rho)^{NF} = -\frac{\sqrt{2}}{3} h_\pi^{NF} \quad (31)$$

follows.

According to Eq. (I.41) and[15] the NF parts of the amplitudes h_π and $A(\Lambda^\circ_-, \Xi^-_-)$ in the MIT bag model have the following structure*

$$h_\pi^{NF} = 4\sqrt{2}\, G(c_1^5 - c_1^6)(I_a - \frac{1}{3} I_b)/f_\pi, \quad (32)$$

$$A(\Lambda^\circ_-)^{NF} = -\sqrt{12}\, G[c^A_{1/2}\, 2(I_a + I_b)/f_\pi + (c^5_{1/2} - c^6_{1/2})(I_a + I_b)/f_\pi], \quad (33)$$

$$A(\Xi^-_-)^{NF} = -4\sqrt{12}\, G[c^A_{1/2}\, (I_a + I_b)/f_\pi + (c^5_{1/2} - c^6_{1/2})\frac{2}{3} I_b/f_\pi], \quad (34)$$

where the overlapping integrals of the quark wave functions of the MIT bag model $I_a = I_a[G^2, F^2]$ and $I_b = I_b[G^2, F^2]$ are determined after Eq. (I.41). This structure remains valid in any other model of independent quarks, only values of I_a and I_b being changed. With increasing mass of quarks the ratio I_b/I_a decreases** and in the nonrelativistic limit it tends to zero.

From comparison of (28)-(30) with (32)-(34) it is seen that the representation (22) reproduces exactly the functional structure of $A^{NF}_{\pi B', B}$ in the nonrelativistic limit ($I_b = 0$) with $b/\sqrt{6} = I_a/f_\pi$. If $I_b \neq 0$ we have

$$b/\sqrt{6} = (I_a + I_b)/f_\pi, \quad (35)$$

while the deviation of (28)-(30) from (32)-(34) appears only in the matrix elements of the mixed operators $\mathcal{O}^{5,6}$. Therefore for the calculation of $(h^\circ_\rho)^{NF}$ we make use of b given by the formula (35) whereas for $(h^1_\rho)^{NF}$ we keep the formula (31). In this way, we arrive at the following expressions for the constants

$$(h^\circ_\rho)^{NF} = -16G(\frac{1}{\sqrt{3}} c_0^A + c^{1A})(I_a + I_b)/f_\pi, \quad (36)$$

$$(h^1_\rho)^{NF} = -\frac{8}{3} G(c_1^5 - c_1^6)(I_a - \frac{1}{3} I_b)/f_\pi, \quad (37)$$

$$(h^2_\rho)^{NF} = (h^\circ_\omega)^{NF} = (h^1_\omega)^{NF} = 0. \quad (38)$$

*In the limit of SU(3) symmetry $m_u = m_d = m_s = 0$.

**Of course the normalization $I_a[G, F] + I_b[G, F] = 1$ is conserved.

4. RESULTS

So, we have found all the contributions to the constants h_V. According to (9) the total values of these constants are determined by the sum $h_V = h_V^F + h_V^{NF}$. Their numerical values for the coefficient functions c_i^R obtained in I, sect.2 are presented in Table 1. In the last column for comparison the "best values" of h_V /9/ are listed. As is seen from the table, our values of h_V agree with $h_V^{b.v.}$. The negative sign of the constant h_ρ^0 is due to its NF part, $(h_\rho^0)^{NF}$, that testifies the arguments given in /9/. Note that h_ρ^0 differs from its bare value $(h_\rho^0)_{\alpha_s=0}$ by a factor of ~ 15, and h_ρ^1 and h_ω^0 even change their signes when quark-gluon interactions switched on.

Table 1

Constants h_V in the model $SU(2)_L \times U(1) \times SU(3)_c$. In brackets their values at $\alpha_s = 0$ are given

	$h_V^F \times 10^7$	$h_V^{NF} \times 10^7$	$h_V \times 10^7$	$h_V^{b.v.} \times 10^7$
h_ρ^0	12.9 (8.3)	-21.2 (-8.8)	-8.3 (-0.55)	-11.4
h_ρ^1	0.24 (-0.15)	0.15 (-0.11)	0.39 (-0.26)	-0.19
h_ρ^2	-6.7 (-11.1)	0	-6.7 (-11.1)	-9.5
h_ω^0	-3.9 (+2.2)	0	-3.9 (+2.2)	-1.9
h_ω^1	-2.2 (-2.2)	0	-2.2 (-2.2)	-1.1

It is interesting to find partial contributions to h_V from charged and neutral currents (Table 2). Here h_V (C.C.) is the parts of h_V defined by charged currents; whereas h_V (N.C.), by neutral currents (see formula (I.4)).

From Table 2 it is seen that the neutral currents, in fact completely, determine the constants h_V^1 (like h_π, see (I.44)), increase by a factor of $1.3\,|h_V^0|$ and decrease by a factor of $2.3\,|h_\rho^2|$. As we shall show in part III of our work the values obtained for h_V testify to the standard model $SU(2)_L \times U(1) \times SU(3)_c$.

Table 2

Contributions of charged (h_V(C.C.)) and neutral (h_V(N.C.)) currents to h_V

	$h_V(C.C.) \times 10^7$	$h_V(N.C.) \times 10^7$	$h_V \times 10^7$
h_ρ^0	- 6.2	- 2.1	- 8.3
h_ρ^1	0.00	0.39	0.39
h_ρ^2	-15.5	8.8	- 6.7
h_ω^0	- 2.9	- 1.0	- 3.9
h_ω^1	0.0	- 2.2	- 2.2

To conclude, we will summarize main points and results of our calculation of h_π and h_V.

The starting point of our scheme is the choice of a low point of renormalization of the operators of the effective Hamiltonian $\mathcal{H}^{PNC}$ (see I, sect. 2). As a result, (in the logarithmic approximation) hadron matrix elements of these operators are determined only by valence quarks, whereas the contributions from all loops, including those of the sea quarks, are taken into account in the coefficient functions of $\mathcal{H}^{PNC}$. Then PNC MNN vertices are represented by sums of contributions of two types different in the quark structure: $h_M = h_M^F + h_M^{NF}$.

The calculation of the F parts of h_π and h_V does not require any model of confinement. Here, of the most interest is h_π^F, since it is proportional to such a fundamental quantity of the theory as $\langle 0|\bar{q}q|0\rangle$. The F parts of h_V are determined by the dominance of vector mesons in vector currents and are expressed only in terms of the experimentally known constants.

To calculate NF parts of h_M one should apply to a certain model of confinement (we make use of the MIT bag model). Unlike h_π^{NF} the NF parts of h_V cannot be calculated directly. For their calculation it is crucial that the matrix elements $\langle VN'|\mathcal{H}^{PNC}|N\rangle^{NF}$ and $\langle \pi B'|\mathcal{H}^{PNC}|B\rangle^{NF}$ have the same quark structure. It is just this fact that has allowed us, by using an approximate SU(6) symmetry of these matrix elements, to express h_V^{NF}, like h_π^{NF}, through the quark wave functions of the MIT bag model. The NF contributions are the most important in h_ρ^0, the NF contributions to the constants h_ρ^2 and $h_\omega^{0,1}$ turn out to be zero.

Our calculation of the constants h_π and h_V does not contain arbitrary (fitting) parameters and artificial assumptions and is self-consistent.

Experimental consequences to which the obtained set of the constants h_M leads will be discussed in part III.

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APPENDIX. Components of the matrix elements (23), (24).

We shall present here results of the calculation of the matrix elements

$$A_M(\mathcal{O}_q) = b < (q\bar{q})_M\,(qqq)_{B'}\,|\mathcal{O}_q|\,(qqq)_B >^{NF}\,, \qquad (A.1)$$

$$A_{\bar{M}}(\mathcal{O}_q) = b\,\eta_C(M) < (qqq)_{B'}\,|\mathcal{O}_q|\,(qqq)_B\; C\,(q\bar{q})_M >^{NF}, \qquad (A.2)$$

$$\mathcal{O}_q = \begin{cases} \bar{q}_a q^b \bar{q}_c q^d = :\bar{q}_{ia}\gamma_\mu\gamma_5 q^{ib}\,\bar{q}_{jc}\gamma^\mu q^{jd}:, \\ (\bar{q}_a q^b \bar{q}_c q^d)' = :\bar{q}_{ia}\gamma_\mu\gamma_5 q^{jb}\,\bar{q}_{jc}\gamma^\mu q^{id}: \end{cases} \qquad (A.3)$$

(see (22)) corresponsing to the amplitudes (25)-(30). In (A1)-(A3) $a, b, c, d = u, d, s$, a sum over colour indices $i, j = 1,2,3$ is assumed. The states $|(q\bar{q})_M>$ and $|(qqq)_{B'}>$ have the following form

$$|(q\bar{q})_M> = \delta^i_j M^A_B b^+_{iA} d^{+jB}|0>, \quad M^{+A}_{\;B} M^B_{\;A} = 1\,, \qquad (A.4)$$

$$|(qqq)_B> = \epsilon^{ijk} B^{ABC} b^+_{iA} b^+_{jB} b^+_{kC}|0>, \quad B^+_{ABC}B^{ABC} = 1\,, \qquad (A.5)$$

$$C\,|(q\bar{q})_M> = \delta^i_j \bar{M}^A_B b^+_{iA} d^{+jB}|0>, \qquad (A.6)$$

where $b^+_{iA}(d^{+iA})$ are operators of creation of a quark (an anti-quark) with the momentum $\vec{p} = 0$; $A, B, C = 1,2,\dots,6$ are SU(6) indices: $1 = u_\uparrow(\bar{u}_\uparrow)$, $2 = u_\downarrow(\bar{u}_\downarrow)$, ..., $6 = s_\downarrow(\bar{s}_\downarrow)$; B^{ABC} are symmetric in A, B, C. The phases of baryon states are taken according to /16/. The phases of $|\pi^->$ and $|V^o_{(o)}>$ are fixed by the matrix elements (see (14) and (I.14)):

$$<0|\frac{1}{2}[(V^3)^1_1 \mp (V^3)^2_2]|\,(\begin{smallmatrix}\rho^o\\ \omega\end{smallmatrix})_{(o)}> = \frac{1}{\sqrt{(2\pi)^3 2m_\rho}}\,\frac{m^2_\rho}{f_\rho}\,, \qquad (A.7)$$

$$<0|(A^o)^2_1|\pi^-> = -\frac{1}{\sqrt{(2\pi)^3 2m_\pi}}\,f_\pi m_\pi\,, \qquad (A.8)$$

where $(V^\mu)^b_a = :\bar{q}_{ia}\gamma^\mu q^{ib}:$; $(A^\mu)^b_a = :\bar{q}_{ia}\gamma^\mu\gamma_5 q^{ib}:$. Now we shall list nonzero components of the functions B^{ABC} of the baryons $p_\uparrow$, $n_\uparrow$, $\Lambda_\uparrow$, $\Xi^-_\uparrow$ and functions M^A_B of the mesons $\rho^o_{(o)}$, $\omega_{(o)}$, π^-:

$$\begin{aligned} p_\uparrow&: \quad B^{114} = -2B^{123} = \frac{1}{9\sqrt{2}}\,, \\ n_\uparrow&: \quad B^{233} = -2B^{134} = -\frac{1}{9\sqrt{2}}\,, \\ \Lambda_\uparrow&: \quad B^{145} = -B^{235} = \frac{1}{12\sqrt{3}}\,, \\ \Xi^-_\uparrow&: \quad B^{455} = -2B^{356} = \frac{1}{9\sqrt{2}}\,; \end{aligned} \qquad (A.9)$$

$$\begin{aligned} \rho^o_{(o)}&: \quad M^1_2 = M^2_1 = -M^3_4 = -M^4_3 = \frac{1}{2\sqrt{3}}\,, \\ &\quad [\bar{M}^1_2 = \bar{M}^2_1 = -\bar{M}^3_4 = -\bar{M}^4_3 = -\frac{1}{2\sqrt{3}}\,, \quad \eta_C(\rho) = -1], \\ \omega_{(o)}&: \quad M^1_2 = M^2_1 = M^3_4 = M^4_3 = \frac{1}{2\sqrt{3}}\,, \\ &\quad [\bar{M}^1_2 = \bar{M}^2_1 = \bar{M}^3_4 = \bar{M}^4_3 = -\frac{1}{2\sqrt{3}}\,, \quad \eta_C(\omega) = -1], \\ \pi^-&: \quad M^3_2 = -M^4_1 = \frac{i}{\sqrt{6}}\,, \\ &\quad [\bar{M}^2_3 = -\bar{M}^1_4 = -\frac{i}{\sqrt{6}}\,, \quad \eta_C(\pi) = 1]; \end{aligned} \qquad (A.10)$$

the isotriplets $\{\rho\}$ and $\{\pi\}$ have the form

$$\{\rho\} = \begin{pmatrix} -\rho^+ \\ \rho^o \\ \rho^- \end{pmatrix}, \qquad \{\pi\} = \begin{pmatrix} -\pi^+ \\ \pi^o \\ \pi^- \end{pmatrix} \qquad (A.11)$$

Contributions of the operators $\bar{q}_a q^b \bar{q}_c q^d$ to the amplitudes (25)-(30) are presented in Tables A1, A2. Matrix elements of the operators $(\bar{q}_a q^b \bar{q}_c q^d)'$, due to the antisymmetry of the state $|(qqq)_B>$ in colour indices, obey the following relation

$$A_{M,\bar{M}}(\bar{q}_a q^b \bar{q}_c q^d)' = -A_{M,\bar{M}}(\bar{q}_a q^b \bar{q}_c q^d)\,. \qquad (A.12)$$

Then it follows that the matrix elements of the symmetric operators $\mathcal{O}^S \propto \bar{q}_a q^b \bar{q}_c q^d + (\bar{q}_a q^b \bar{q}_c q^d)'$ vanish (see /14/).

Table A1.

Contributions of the matrix elements $A_{M,\bar{M}}(O_q)$ to the amplitudes (23), (24) with $\Delta S = 0$.

		$\bar{u}u\bar{u}u$	$\bar{d}d\bar{d}d$	$\bar{u}u\bar{d}d$	$\bar{d}d\bar{u}u$	$\bar{u}d\bar{d}u$	$\bar{d}u\bar{u}d$
$\langle\rho^0_{(\omega)} p_\uparrow\vert O_q\vert p_\uparrow\rangle^{NF}$	A_M	0	0	$\frac{4}{3\sqrt{3}}b$	$\frac{2}{3\sqrt{3}}b$	$-\frac{4}{3\sqrt{3}}b$	$-\frac{2}{3\sqrt{3}}b$
	$A_{\bar{M}}$	0	0	$\frac{4}{3\sqrt{3}}b$	$\frac{2}{3\sqrt{3}}b$	$-\frac{2}{3\sqrt{3}}b$	$-\frac{4}{3\sqrt{3}}b$
	$A_M+A_{\bar{M}}$	0	0	$\frac{8}{3\sqrt{3}}b$	$\frac{4}{3\sqrt{3}}b$	$-\frac{2}{\sqrt{3}}b$	$-\frac{2}{\sqrt{3}}b$
$\langle\omega_{(0)} p_\uparrow\vert O_q\vert p_\uparrow\rangle^{NF}$	A_M	0	0	0	0	$-\frac{1}{\sqrt{3}}b$	$-\frac{1}{\sqrt{3}}b$
	$A_{\bar{M}}$	0	0	0	0	$-\frac{1}{\sqrt{3}}b$	$-\frac{1}{\sqrt{3}}b$
	$A_M+A_{\bar{M}}$	0	0	0	0	0	0
$\langle\pi^- p_\uparrow\vert O_q\vert n_\uparrow\rangle^{NF}$	A_M	0	0	$-\frac{i}{\sqrt{6}}b$	$\frac{i}{\sqrt{6}}b$	$\frac{i}{\sqrt{6}}b$	$-\frac{i}{\sqrt{6}}b$
	$A_{\bar{M}}$	0	0	$-\frac{i}{\sqrt{6}}b$	$\frac{i}{\sqrt{6}}b$	$-\frac{i}{\sqrt{6}}b$	$\frac{i}{\sqrt{6}}b$
	$A_M+A_{\bar{M}}$	0	0	$-\frac{2i}{\sqrt{6}}b$	$\frac{2i}{\sqrt{6}}b$	0	0

Table A2.

Contributions of the matrix elements $A_{M,\bar{M}}(O_q)$ to the amplitudes (24) with $\Delta S = 1$ (continued on the next page)

		$\bar{u}u\bar{d}s$	$\bar{d}s\bar{u}u$	$\bar{u}d\bar{s}u$	$\bar{s}u\bar{u}d$	$\bar{u}s\bar{d}u$
$\langle\pi^- p_\uparrow\vert O_q\vert\Lambda_\uparrow\rangle^{NF}$	A_M	0	ib	0	0	$-\frac{i}{2}b$
	$A_{\bar{M}}$	$-\frac{i}{2}b$	$\frac{i}{2}b$	0	0	$-\frac{i}{2}b$
	$A_M+A_{\bar{M}}$	$-\frac{i}{2}b$	$\frac{3i}{2}b$	0	0	$-ib$
$\langle\pi^-\Lambda_\uparrow\vert O_q\vert\Xi^-_\uparrow\rangle^{NF}$	A_M	$\frac{i}{2}b$	$\frac{3i}{2}b$	0	0	$-\frac{3i}{2}b$
	$A_{\bar{M}}$	0	0	0	0	0
	$A_M+A_{\bar{M}}$	$\frac{i}{2}b$	$\frac{3i}{2}b$	0	0	$-\frac{3i}{2}b$

		$\bar{d}u\bar{u}s$	$\bar{d}d\bar{d}s$	$\bar{d}s\bar{d}d$	$\bar{d}s\bar{s}s$	$\bar{s}sds$
$\langle\pi^- p_\uparrow\vert O_q\vert\Lambda_\uparrow\rangle^{NF}$	A_M	$-\frac{i}{2}b$	0	0	0	0
	$A_{\bar{M}}$	$\frac{i}{2}b$	$\frac{i}{2}b$	$-\frac{i}{2}b$	0	0
	$A_M+A_{\bar{M}}$	0	$\frac{i}{2}b$	$-\frac{i}{2}b$	0	0
$\langle\pi^-\Lambda_\uparrow\vert O_q\vert\Xi^-_\uparrow\rangle^{NF}$	A_M	$-\frac{i}{2}b$	0	0	0	0
	$A_{\bar{M}}$	0	ib	$-ib$	$\frac{i}{2}b$	$-\frac{i}{2}b$
	$A_M+A_{\bar{M}}$	$-\frac{i}{2}b$	ib	$-ib$	$\frac{i}{2}b$	$-\frac{i}{2}b$

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Дубовик В.М., Зенкин С.В. E2-83-615
Самосогласованный расчет слабых констант в несохраняющих четность ядерных силах. НЧ в ρNN и ωNN вершинах

На основе эффективного несохраняющего четность /НЧ/ гамильтониана стандартной модели $SU(2)_L \times U(1) \times SU(3)_c$, полученного в нашей предыдущей работе, рассмотрена кварковая структура НЧ ρNN и ωNN вершин. Без свободных /подгоночных/ параметров рассчитаны все вклады в константы $h_\rho^{0,1,2}$ и $h_\omega^{0,1}$, входящие в НЧ NN потенциал. Нефакторизуемые /НФ/ вклады в эти константы рассчитаны, исходя из приближенной SU(6) симметрии матричных элементов $\langle MB'|\mathcal{H}^{НЧ}|B\rangle^{НФ}$, в модели массачусетского мешка. Полученные значения $h_{\rho,\omega}$ находятся в согласии с "лучшими значениями" $h_{\rho,\omega}$ Desplanques, Donoghue и Holstein.

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Dubovik V.M., Zenkin S.V. E2-83-615
Self-Consistent Calculation of the Weak Constants in the Parity Nonconserving Nuclear Forces. PNC in the ρNN and ωNN Vertices

The quark structure of the parity nonconserving (PNC) ρNN and ωNN vertices is considered for the effective PNC Hamiltonian of the standard model $SU(2)_L \times U(1) \times SU(3)_c$ obtained in our previous paper. Without arbitrary (fitting) parameters all the contributions to the constants $h_\rho^{0,1,2}$ and $h_\omega^{0,1}$ of the PNC NN potential are calculated. Nonfactorizable (NF) contributions to the constants are calculated proceeding from the approximate SU(6) symmetry of matrix elements $\langle MB'|\mathcal{H}^{PNC}|B\rangle^{NF}$, in the MIT bag model. The obtained values of $h_\rho^{0,1,2}$ and $h_\omega^{0,1}$ agree with the "best values" of $h_{\rho,\omega}$ of Desplanques, Donoghue and Holstein.

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

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