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WEAK N-N INTERACTION IN THE SIMPLEST REACTIONS.

I. State of the Problem. General Form of the Long-Range Parity-Violating Potential

Дубовик В.М., Замиралов В.С., Зенкин С.В. **E2 - 12381**

Слабые NN-взаимодействия в простейших реакциях.
1. Состояние проблемы. Общий вид дальнодействующего потенциала, не сохраняющего четность

На основе токов SU(4) модели Вайнберга-Салама рассмотрены нарушающие четность ρ, ω - обменные потенциалы. В этих рамках пересмотрены и уточнены наиболее распространенные параметризации потенциалов. Оценены относительные вклады сепарабельных и несепарабельных диаграмм в параметр асимметрии A_{pp} в реакции $\vec{p}+p \to p+p$ и циркулярную поляризацию γ -квантов в реакциях $n+p \to d+\gamma$ и $n+d \to t+\gamma$. Рассмотрение проведено в различных модификациях приближения факторизации и для различных сильных потенциалов. Обращено внимание на отсутствие вклада ω-мезона в сепарабельную часть P_γ в реакции $n+p \to d+\gamma$. Несепарабельные вклады оценены в простой калибровочной модели.

Работа выполнена в Лаборатории теоретической физики ОИЯИ.

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Dubovik V.M., Zamiralov V.S., Zenkin S.V. **E2- 12381**

Weak N-N Interactions in the Simplest Reactions.
1. State of the Problem. General Form of the Long-Range Parity-Violating Potential

Parity-violating ρ,ω-exchange potentials based on SU(4) currents of the Weinberg-Salam model are considered. The conventional parametrizations of weak potentials are revised and specified within SU(4) currents. The relative contributions of separable and nonseparable diagrams to the asymmetry parameter A_{pp} in the reaction $\vec{p}+p \to p+p$ and circular γ-polarization P_γ in the reactions $n+p \to d+\gamma$ and $n+d \to t+\gamma$ are evaluated within the different modifications of the factorization approach and various strong potentials. The absence of ω-meson contribution to the separable part of the P_γ in the reaction $n+p \to d+\gamma$ is regarded. Nonseparable contributions are also evaluated within the simple gauge model.

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

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1. INTRODUCTION

The Weinberg-Salam model of the unified theory of weak and electromagnetic interactions /1/ successfully describes lepton-lepton and lepton-hadron interactions with the value of $\sin^2\theta_W$ equal to ~0.25 (ref. /2/). At the same time the hadron-hadron weak interaction imposes many problems /3/. The most serious one is that NN-weak interaction effects yield an unexpectedly large experimental value of the circular γ-polarization $P_\gamma = (-1.30 \pm 0.45) \times 10^{-6}$ in the process $n+p \to d+\gamma$ (ref. /4/), which exceeds by 10-100 times the theoretical predictions (see reviews /5/). The study of the weak NN-interaction is now even more actual, as high-precision experiments are planned to measure the weak correlation effects in the processes $n+p \to d+\gamma$, $\gamma+d \to n+p$, $e+d \to e+n+p$, as well as the spin precession of the cold neutrons in matter. These experiments would help to clear the source of discrepancy, which could be either experimental errors due to extraordinary measurement difficulties and/or our insufficient knowledge of the weak-interaction dynamics in the hadron-hadron collisions. Moreover strong discrepancy can arise also due to some peculiar behaviour of strong potential or even due to some unusual phase state of hadronic matter at the very short distances characteristic of the weak interaction.

Calculations of such complex effects forced to make many assumptions not only of principal but also of technical character. In the earlier papers the diagram of Fig. 1 was supposed to dominate in the weak NN-interaction, and the main hypothesis were the validity of the current-field identity and the possibility of saturating matrix elements with vacuum intermediate

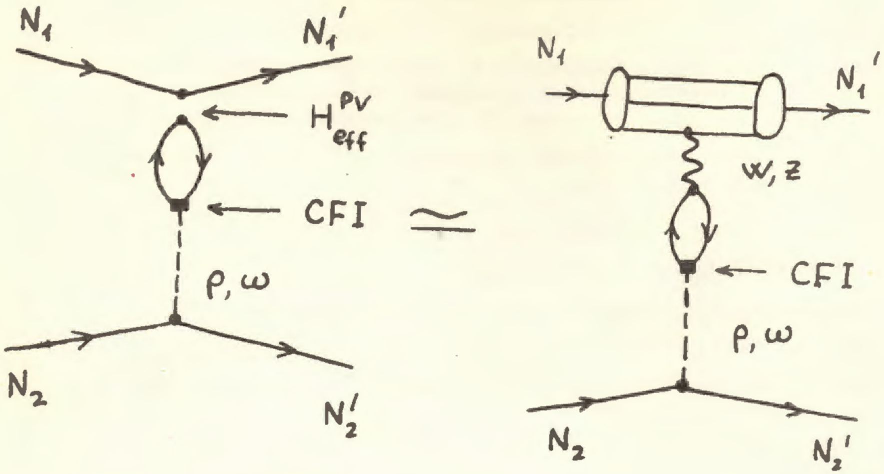

Fig. 1

states only /6/. However, for various models of weak interactions and for various local nuclear potentials it was proved to be impossible to reach accord with the experimental value P_γ^{exp}. Neither it was possible to reach a consistent description of a series of experiments on violation of parity in the hadron-hadron interaction /7/.

The next step was to take into account not only "separable" diagrams of Fig. 1, but also "non-separable" diagram of Fig. 2. The calculation of these latter diagrams required new important assumptions for solving a complicated problem of exchange between nucleon lines of the bound quark-antiquark states. (In the last years the weak interaction is considered in the NN -collisions on the level of quarks rather than on the level of nucleons). In spite of great effects, the contribution of the nonseparable diagrams at most doubled the result for P_γ^{th} /8,9/ .

The last step was to take into account the gluon corrections to the weak interaction vertices within QCD /8,9/. However, in practice the theoretical value

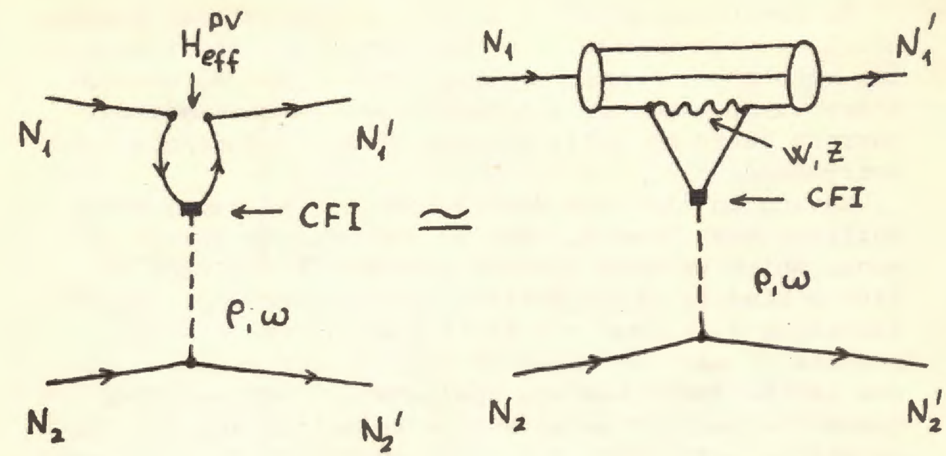

Fig. 2

of P_γ was either enhanced by a factor of ~ 2.5 /9/ or even suppressed by a factor of ~ 4 /8/. The discrepancy between these results seems to lie in different ways of treating non-separable contributions.

One can conclude that the value of P_γ^{th} has proved to be remarkably stable to all enhancement mechanisms, giving essentially the result close to that predicted by Cabibbo charged current theory, that is $P_\gamma \sim (2-3)\times 10^{-8}$. But at this stage the number of assumptions is great, and it is difficult to understand whether the value of P_γ^{th} is intrinsically small or some peculiar cancellations are present. Moreover, some hypothesis may even be inconsistent or lead to double-counting. This can occur, for example, if one works with two different models of strong interaction, with coloured massless gluons and massive vector bosons, at the same time. We think that only unambiguous construction of the strong exchange potentials in the framework of QCD valid for all ranges, would justify such a procedure.

By now there exists a number of theoretical schemes which however differ from each other by definitions of the constants, normalizations, and so on. Moreover, there is no explicit expression of the weak neutral current based on SU(4) instead of SU(3) currents used everywhere.

Having in mind the development of the gauge model unifying weak bosons, photons and massive vector mesons, which we have started earlier /10/, we want to give a kind of short review of the preceding results (sections 2,3). Our aim is to present the essential results of various works on the subject obtained by now in the factorization approach and its modifications in order to be able to make selfconsistent calculations and a detailed comparison with the results of other works. We consider here polarization of the photons in the process $n+p\to d+\gamma$ and also in $n+d\to t+\gamma$, and asymmetry parameter in the elastic pp -scattering with the polarized proton beam, disregarding for a time those processes where π -meson exchange is essential.

2. GENERAL FORM OF H_W AND ANALYSIS OF SOME PRECEDING RESULTS IN A SIMPLE FACTORIZATION APPROACH

Here, not pretending to analyse all the works on the subject, we shall try to consider the main lines along which NN -interaction was studied. Besides, although most of the authors work with different models of weak interaction we take only that of Weinberg and Salam /1/ within the four-quark GIM scheme /11/. A strong NN -interaction is described as usual by a certain vector-boson-exchange potential as shown in Fig. 3. In order to take into account the NN weak interaction effects, having in mind the long-range nuclear forces due to the existence of strong core, it is necessary to introduce the weak interaction into one of the vertices of Fig. 3. As the weak NN -interactions are observed only due to the

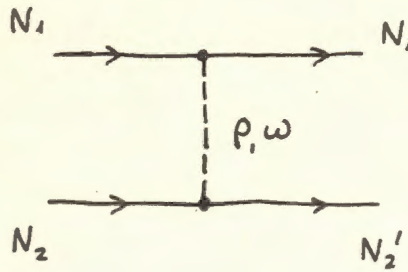

Fig. 3

parity-violating effects, the central point is to construct parity-violating nucleon-vector-meson vertex. The most ingenious way to do it is to consider the so-called "separable" contribution given by the diagram of Fig. 1, where H_{eff}^{PV} means a parity-violating part of the effective current-current weak Hamiltonian

$$H_{eff} = \frac{G}{\sqrt{2}}\left(J_\lambda^{C+} J_\lambda^{C} + J_\lambda^{N} J_\lambda^{N}\right), \qquad (1)$$

where G is the Fermi constant*.

The charged weak current in the GIM model reads

$$J_\lambda^C = [\,\bar d\,\gamma_\lambda(1+\gamma_5)\,u + \bar s\,\gamma_\lambda(1+\gamma_5)\,c\,]\cos\theta_C + [\,-\bar d\,\gamma_\lambda(1+\gamma_5)\,c + \bar s\,\gamma_\lambda(1+\gamma_5)\,u\,]\sin\theta_C\,, \qquad (2)$$

where θ is the Cabibbo angle, $\sin^2\theta_C = 0.055$, while neutral weak current is put in the form

$$J_\lambda^N = \sum_{k=0,3,8,15} \left(v^k V_\lambda^k + a^k A_\lambda^k \right), \qquad (3)$$

where

$$V_\lambda^0 = \frac{1}{2\sqrt{2}}\left(\bar c\gamma_\lambda c + \bar u\gamma_\lambda u + \bar d\gamma_\lambda d + \bar s\gamma_\lambda s\right), \quad V_\lambda^3 = \frac{1}{2}\left(\bar u\gamma_\lambda u - \bar d\gamma_\lambda d\right),$$

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* We use Pauli metric.

$$V^8_\lambda = \frac{1}{2\sqrt{3}}(\bar{u}\gamma_\lambda u + \bar{d}\gamma_\lambda d - 2\bar{s}\gamma_\lambda s),\quad V^{15}_\lambda = \frac{1}{2\sqrt{6}}(-3\bar{c}\gamma_\lambda c + \bar{u}\gamma_\lambda u + \bar{d}\gamma_\lambda d + \bar{s}\gamma_\lambda s), \tag{4}$$

analogous expressions being valid for axial-vector currents with $\gamma_\lambda \to \gamma_\lambda\gamma_5$. In the Weinberg-Salam scheme with four quarks, the neutral current (3) reduces to the expression

$$\begin{aligned} J^N_\lambda &= \left(A^3_\lambda - \frac{2}{\sqrt{6}}A^{15}_\lambda + \frac{1}{\sqrt{3}}A^8_\lambda\right) - \frac{2\sqrt{2}}{3}\sin^2\theta_W V^0_\lambda + \\ &+ (1 - 2\sin^2\theta_W)\left(V^3_\lambda - \frac{2}{\sqrt{6}}V^{15}_\lambda + \frac{1}{\sqrt{3}}V^8_\lambda\right), \end{aligned} \tag{5}$$

where θ_W is the Weinberg angle; so the coefficients v^k, a^k of Eq. (3) are

$$\begin{aligned} v^0 &= -\frac{2\sqrt{2}}{3}\sin^2\theta_W, \qquad a^0 = 0, \\ v^3 &= \sqrt{3}\,v^8 = -\frac{\sqrt{6}}{2}v^{15} = 1 - 2\sin^2\theta_W, \\ a^3 &= \sqrt{3}\,a^8 = -\frac{\sqrt{6}}{2}a^{15} = 1. \end{aligned} \tag{6}$$

If we now retain only u, d quarks, as it is reasonable in the case of NN-interaction, the sum $J^8_\lambda - \sqrt{2}J^{15}_\lambda$ disappears. Hence,

$$\begin{aligned} J^N_\lambda &= A^3_\lambda + (1 - 2\sin^2\theta_W)V^3_\lambda - \frac{2}{3}\sin^2\theta_W \tilde{V}^0_\lambda, \\ \tilde{V}^0_\lambda &= \frac{1}{2}(\bar{u}\gamma_\lambda u + \bar{d}\gamma_\lambda d). \end{aligned} \tag{7}$$

Now let us write a general form of the nonstrange parity-violating part of the H_{eff}(1) in the framework of the SU(4) scheme:

$$\begin{aligned} H^{PV}_{eff} = \frac{G}{\sqrt{2}}\Big[& A\left(V^{1+i2}_\lambda A^{1-i2}_\lambda + V^{1-i2}_\lambda A^{1+i2}_\lambda\right) + B V^3_\lambda A^3_\lambda + \\ & + C V^3_\lambda A^8_\lambda + C' V^8_\lambda A^3_\lambda + D V^8_\lambda A^8_\lambda + E V^0_\lambda A^0_\lambda + F V^3_\lambda A^0_\lambda + F' V^0_\lambda A^3_\lambda + \\ & + G V^8_\lambda A^0_\lambda + G' V^0_\lambda A^8_\lambda + K V^{15}_\lambda A^{15}_\lambda + P V^3_\lambda A^{15}_\lambda + P' V^{15}_\lambda A^3_\lambda + \\ & + M V^8_\lambda A^{15}_\lambda + M' V^{15}_\lambda A^8_\lambda + L V^{15}_\lambda A^0_\lambda + L' V^0_\lambda A^{15}_\lambda \Big]. \end{aligned} \tag{8}$$

One can easily read off the coefficients of Eq. (8) using Eqs. (3-6). For the Weinberg-Salam model

$$\begin{aligned} & A = \cos^2\theta_C, \\ & B = 3D = \sqrt{3}\,C = \sqrt{3}\,C' = \frac{3}{2}K = 2(1 - 2\sin^2\theta_W), \\ & \frac{\sqrt{6}}{2}P = \frac{\sqrt{6}}{2}P' = \frac{3}{\sqrt{2}}M = \frac{3}{\sqrt{2}}M' = -2(1 - 2\sin^2\theta_W), \\ & F' = \sqrt{3}\,G' = -\sqrt{\frac{3}{2}}L' = -\frac{2\sqrt{2}}{3}2\sin^2\theta_W, \\ & E = F = G = L = 0. \end{aligned} \tag{9}$$

Eq. (8) is reduced to that of ref. /9/ if we omit c-quark altogether, that is if we put $J^0_\lambda = \sqrt{3}\,J^{15}_\lambda$ and take into account the difference in normalization by $J^0_\lambda = \frac{\sqrt{3}}{2}(J^0_\lambda)^D$, where index D means "taken from /9/ ". The coefficients of ref. /9/ are expressed through those of Eq. (9) as follows:

$$\begin{aligned} & A^D = A,\quad B^D = B,\quad C^D = C,\quad C'^D = C',\quad D^D = D, \\ & E^D = \frac{3}{4}\left(E + \frac{1}{3}K + \frac{L + L'}{\sqrt{3}}\right), \\ & F^D = \frac{1}{2}(\sqrt{3}F + P),\quad F'^D = \frac{1}{2}(\sqrt{3}F' + P'), \\ & G^D = \frac{1}{2}(\sqrt{3}G + M),\quad G'^D = \frac{1}{2}(\sqrt{3}G' + M'). \end{aligned} \tag{10}$$

The next step of the factorization approach is to calculate the matrix elements of H^{PV}_{eff} between the vector meson states and $N\bar{N}'$ state which yields

$$\langle \rho/\omega \,|\, V^i_\lambda A^j_\lambda \,|\, N\bar{N}'\rangle = \sum_n \langle \rho/\omega \,|\, V^i_\lambda \,|\, n\rangle\langle n \,|\, A^j_\lambda \,|\, N\bar{N}'\rangle \simeq$$
$$\simeq \langle \rho/\omega \,|\, V^i_\lambda \,|\, 0\rangle\langle 0 \,|\, A^j_\lambda \,|\, N\bar{N}'\rangle \,, \quad i,j = 0,1,2,3,8,15\,, \tag{11}$$

where it is assumed as usual that it is possible to retain only vacuum intermediate state [5]. The matrix element of A^3_λ is just g_A, while that of A^8_λ can be found through the SU(3) relation

$$\xi = \frac{1}{2}\,\frac{\langle 0 \,|\, A^8_\lambda \,|N\bar{N}'\rangle}{\langle 0 \,|\, A^3_\lambda \,|N\bar{N}'\rangle} = -\frac{1}{2\sqrt{3}}\,\frac{1-\dfrac{3F}{D}}{1+\dfrac{E}{D}} \tag{12}$$

and $\xi^{-1} = 6\sqrt{3}$ with SU(3) F/D ratio equal to 1/2 as taken in refs. [6,8] and $\xi = \sqrt{3}/10$ with F/D = 2/3 as taken in ref. [9]. Neglecting charm quarks, we can define also

$$x = \frac{\sqrt{2}}{9}\,\frac{\langle 0|A^0_\lambda \,|\, N\bar{N}'\rangle}{\langle 0|\, A^8_\lambda \,|\, N\bar{N}'\rangle} = \frac{\sqrt{2}}{3\sqrt{3}}\,\frac{\langle 0 \,|\, A^{15}_\lambda \,|N\bar{N}'\rangle}{\langle 0 \,|\, A^8_\lambda \,|\, N\bar{N}'\rangle} = \frac{1}{3\sqrt{6}}\,\frac{\langle 0 \,|\, A^{0D}_\lambda \,|N\bar{N}'\rangle}{\langle 0 \,|\, A^8_\lambda \,|\, N\bar{N}'\rangle} \tag{13}$$

and for $x^{-1} = 3\sqrt{3}$, one obtains the result quoted in ref. [9], while in [8] $x = \sqrt{3}/5$ was used.

Similarly, neglecting c-quark, the following expression was defined in ref. [8]

$$y = \sqrt{2}\,\frac{\langle 0|\, V^0_\lambda \,|\, N\bar{N}'\rangle}{\langle 0|\, V^8_\lambda|\, N\bar{N}'\rangle} = \sqrt{6}\,\frac{\langle 0 \,|\, V^{15}_\lambda |N\bar{N}'\rangle}{\langle 0 \,|\, V^8_\lambda |N\bar{N}'\rangle} = \sqrt{\frac{3}{2}}\,\frac{\langle 0 \,|\, V^{0D}_\lambda |N\bar{N}'\rangle}{\langle 0 \,|\, V^8_\lambda \,|\, N\bar{N}'\rangle} \tag{14}$$

and y was taken to be equal to $\sqrt{3}$ in refs. [8,9].

The matrix elements of the vector currents are usually taken from current-field identities, which we choose in the form

$$V^k_\lambda = \frac{m^2_\rho}{f_\rho}\,\rho^k_\lambda\,, \quad \tilde{V}^0_\lambda = \frac{1}{2}(\bar{u}\,\gamma_\lambda u + \bar{d}\,\gamma_\lambda d) = \frac{m^2_\rho}{f_\rho}\,\omega_\lambda\,, \quad k=1,2,3 \tag{15}$$

which is identical with the relations given in ref. [9]

$$\langle \rho|\, J^{em}_\lambda \,|\, 0\rangle = C_\rho m^2_\rho\, \epsilon^*_\lambda\,, \quad \langle \omega \,|\, J^{em}_\lambda |0\rangle = C_\omega\, m^2_\rho\, \epsilon^*_\lambda = \frac{1}{\sqrt{6}}\langle \omega|J^{0D}_\lambda|0\rangle \tag{16}$$

with $C_\rho = 3C_\omega = f^{-1}_\rho$. In refs. [6,8] unphysical ω^8_λ belonging to pure SU(3) octet was apparently used.

The effective Hamiltonian of the parity-violating NNV-vertex can now be put in the standard form

$$H^{PV}_{NNV} = \frac{-i\sqrt{2}Gm^2_\rho g_A}{f_\rho}\Big[\frac{A}{\sqrt{2}}\bar{N}\gamma_\lambda\gamma_5(\tau^+\rho^-_\lambda - \tau^-\rho^+_\lambda)\bar{N} + \frac{B}{4}\bar{N}\gamma_\lambda\gamma_5\tau^3 N\rho^0_\lambda +$$
$$+\frac{\xi}{2}(C + 3\sqrt{6}\,xF)\,\bar{N}\gamma_\lambda\gamma_5 N\rho^0_\lambda + \frac{1}{4}\Big(C' + \sqrt{\frac{2}{3}}\,y\,F'\Big)\bar{N}\gamma_\lambda\gamma_5\tau^3 N\,\frac{\omega_\lambda}{\sqrt{3}} +$$
$$+\frac{\xi}{2}\Big(D + 6xyE + 3\sqrt{6}\,xG + \sqrt{\frac{2}{3}}\,yG'\Big)\bar{N}\gamma_\lambda\gamma_5\bar{N}\,\frac{\omega_\lambda}{\sqrt{3}}\Big]\,, \tag{17}$$

where we have used the coefficients of ref. [9] through Eq. (10) instead of ours to shorten the formula and have omitted the superscript D.

Note, that we return to the form given in [12] disregarding $1/\sqrt{3}$ in terms with ω_λ in Eq. (17), that is effectively using unphysical ω^8_λ.

It is convenient to write here a relation between the coefficients of ref. [9] and those of [8] indicated by index GGT; both for the case of separable contributions:

$$A^D = A^{GGT}\cos^2\theta_C\,, \qquad B^D = 4B^{GGT}\cos^2\theta_C\,,$$

$$C^D = 2(E-M)^{GGT}\cos^2\theta_C\,, \quad C'^D = 2(E+M)^{GGT}\cos^2\theta_C\,,$$

$$F^D = 2(F-K)^{GGT}\cos^2\theta_C\,, \quad F'^D = 2(F+K)^{GGT}\cos^2\theta_C\,,$$

$$D^D = 4C^{GGT}\cos^2\theta_C\,, \qquad E^D = 4D^{GGT}\cos^2\theta_C\,,$$
$$G^D = 2\,(G-L)^{GGT}\cos^2\theta_C\,, \qquad G'^D = 2\,(G+L)^{GGT}\cos^2\theta_C\,. \tag{18}$$

In order to retain the initial form of H^{PV}_{NNV} of ref. [6], as given in ref. [12], we rewrite Eq. (17) in the form

$$H^{PV}_{NNV} = \frac{-i\sqrt{2}Gm_\rho^2 g_A}{f_\rho}\Big[\frac{A}{\sqrt{2}}\bar{N}\gamma_\lambda\gamma_5(\tau^+\rho^-_\lambda - \tau^-\rho^+_\lambda)N + \frac{B}{4}\bar{N}\gamma_\lambda\gamma_5\tau^3 N\rho^0_\lambda + \frac{\xi}{2}\tilde{C}\,\bar{N}\gamma_\lambda\gamma_5 N\rho^0_\lambda + \frac{\tilde{C}'}{4}\bar{N}\gamma_\lambda\gamma_5\tau^3 N\frac{\omega_\lambda}{\sqrt{3}} + \frac{\xi}{2}\tilde{D}\,\bar{N}\gamma_\lambda\gamma_5 N\frac{\omega_\lambda}{\sqrt{3}}\Big]. \tag{19}$$

with obvious redefinitions, and for $x^{-1}=3\sqrt{3}$, $y=\sqrt{3}$ in the Weinberg-Salam model

$$A = \cos^2\theta_C\,, \qquad B = 2\,(1-2\sin^2\theta_W)\,,$$
$$\tilde{C} = C + \sqrt{2}F = 0\,, \quad \tilde{C}' = C' + \sqrt{2}F' = -\frac{4}{\sqrt{3}}\sin^2\theta_W\,, \tag{20}$$
$$\tilde{D} = D + 2E + \sqrt{2}G + \sqrt{2}G' = 0\,.$$

In order to construct the parity-violating weak long-range potential V^{PV}_W based on the diagrams of Fig. 1, it is necessary to write down also strong-interaction effective NNV -Hamiltonian which we write here in the form of ref. [13]

$$H_S = -ig_{\rho NN}\bar{N}\Big[\gamma_\lambda + (\mu_V - 1)\frac{\sigma_{\lambda\mu}}{4M_N}\partial_\mu\Big]\frac{\tau^i}{2}N\rho^i_\lambda - ig_{\omega NN}\bar{N}\Big[\gamma_\lambda + (\mu_S - 1)\frac{\sigma_{\lambda\mu}}{4M_N}\partial_\mu\Big]N\omega_\lambda\,, \tag{21}$$

where μ_V and μ_S are the isovector and isoscalar nucleon magnetic moments, $2i\sigma_{\lambda\mu} = [\gamma_\lambda, \gamma_\mu]_-$, and strong coupling constants were taken in refs. [9,13] as $g^2_{\rho NN}/4\pi = 2{,}4$, $g^2_{\omega NN}/4\pi = 5{,}4$. Instead, the SU(3) relation $g_{\omega NN} = \frac{\sqrt{3}}{2}g_{\rho NN}$ valid for unphysical ω^8_λ was used in refs. [6,12].

Now it is possible to write V^{PV}_W in the form

$$V^{PV}_W = \frac{g_{\rho NN}Gm_\rho^2 g_A}{8\sqrt{2}\pi M_N f_\rho}\Big\{(v+\mu_V w)AT^{(+)} + (v+\mu_V w)B\tau^3_1\tau^3_2 + (v+\mu_S w)\frac{g_{\omega NN}}{\sqrt{3}g_{\rho NN}}\xi\tilde{D} + \Big[(v_1\tau^3_1 - v_2\tau^3_2) + \mu_S w\frac{\tau^3_1+\tau^3_2}{2}\Big]\frac{g_{\omega NN}}{\sqrt{3}g_{\rho NN}}\frac{\tilde{C}'}{2} + \Big[(v_1\tau^3_2 - v_2\tau^3_1) + \mu_V w\frac{\tau^3_1+\tau^3_2}{2}\Big]\frac{\xi\tilde{C}}{2}\Big\}, \tag{22}$$

$$v = v_1 - v_2\,, \quad v_j = \sigma_j\{\vec{p}_{21}, f(r)\}_+\,, \quad w_j = i(\vec{\sigma}_1\times\vec{\sigma}_2)[\vec{p}_{21}, f(r)]_-\,,$$

$$\vec{p}_{21} = \vec{p}_2 - \vec{p}_1\,, \quad f(r) = \frac{1}{r}e^{-m_\rho r}\,, \quad T^{(+)} = \tau^+_1\tau^-_2 + \tau^-_1\tau^+_2\,,$$

$$\tau^\pm_{1,2} = \frac{1}{2}(\tau^1 \pm i\tau^2)_{1,2}\,,$$

and this expression coincides with that given in refs. [6,8] if we take into account the difference between ω_λ and ω^8_λ, that is, if we put instead of $g_{\omega NN}/(\sqrt{3}g_{\rho NN})$ factor $\sqrt{3}/2$ use Eqs. (17-19), and also take $g_{\rho NN} = f_\rho$. Moreover, we obtain the separable part of the expression for V^{PV}_W given in [9] with $C_\rho = f^{-1}_\rho = g^{-1}_{\rho NN} = 3C_\omega$ (see also Eq. (29) below). The diagonal part of Eq. (22), that is, without the terms proportional to $\tilde{C}$, $\tilde{C}'$, is in accord with that given in refs. [13] and [14] and used in ref. [10].

One can observe from Eq. (22) that ω-meson does not contribute to the polarization P_γ in the reaction $n + p \to d + \gamma$ in the approach based only on separable contributions of Fig. 1 within the standard Weinberg-Salam model. Indeed, the ω -contribution is present in P_γ only in the term proportional to $\tilde{D}$, and $\tilde{D}^{WS} = 0$.

The absence of the ω-contribution to P_γ^{WS} can be read off immediately from Eq. (7) and was noted already in refs.[13,15] whereas in ref.[16] it was overseen. In ref.[8] $x=\sqrt{3}/5$, $y=\sqrt{3}$, were used, so ω-contribution persisted even for separable contributions, although small. For $x^{-1}=3\sqrt{3}$ the ω-contribution is absent, and this is essentially the result quoted in ref.[9]. If J_λ^8 is retained in Eq. (5) disregarding J_λ^{15}, one obtains also nonzero ω-contribution to P_γ [12]. This result has been corrected in ref.[17].

We quote now some characteristic values of the polarization P_γ^{th} in the reaction $n+p\to d+\gamma$ and of asymmetry parameter A_{pp}^{th} of the reaction $\vec{p}+p\to p+p$ with polarized proton beam within the outlined approach based on V_W^{PV} and several strong-interaction NN-potentials. Our aim is to illustrate the factorization scheme results within the Weinberg-Salam model.

We have used the expressions for P_γ obtained in refs.[12,13]* with the coefficients of H_{NN}^{PV} given by Eq. (20) at the value $\sin^2\theta_W=0.25$ and

$$f_V = -\frac{\sqrt{2}\,G\,g_A m_\rho^2}{g_{\rho NN}} = -2.127\times 10^{-6}\ , \tag{23}$$

that is,

$$P_\gamma^{HJ} = (-14A+5.7B-0.63\tilde{D})\times 10^{-3} f_V = 1.60\times 10^{-8}$$

for Hamada-Johnston's potential,

$$P_\gamma^{KSW} = (-6.1A+8.2B-0.82\tilde{D})\times 10^{-3} f_V \doteq -0.52\times 10^{-8}$$

for Kishi-Sawada-Watari's one,

$$P_\gamma^{RSC} = (11.3A-6.0B+0.21\tilde{D})\times 10^{-3} f_V = -1.0\times 10^{-8} \tag{24}$$

*We have not succeeded in obtaining the sign of the D-wave contribution cited by ref.[12].

for Reid's soft-core potential, and

$$P_\gamma^{T} = -(10A+2.8B+0.86\tilde{D})\times 10^{-3} f_V = 2.60\times 10^{-8}$$

for Tamagaki's Gaussian soft-core potential.

The asymmetry parameter A_{pp} for $\vec{p}+p\to p+p$ for the Hamada-Johnston potential is found by the formula

$$A_{pp}(E_p=15\,\text{MeV}) = -[2.68\,(B+2\xi\tilde{C})+1.52(\tilde{C}'+2\xi\tilde{D})]\times 10^{-8} = -1.80\times 10^{-8} \tag{25}$$

We have also calculated the polarization parameter P_γ^{tr} for the reaction of n-d capture $n+{}^2H\to{}^3H+\gamma$, using the formulas of ref.[19] with $x=1$:

$$P_\gamma^{tr} = \kappa\,(7.224A+0.545B-0.0069\tilde{C}+0.2236\tilde{C}'-0.0734\tilde{D}) = 1.21\times 10^{-6}, \tag{26}$$

where $\kappa = 1.67\times 10^{-7}$. Note that this result obtained with $\sin^2\theta_W=0.25$ is practically the same as that of ref.[19] obtained with Cabibbo currents only. (We have also used the fact that π-meson contribution is negligibly small[19]).

3. MODIFICATIONS OF A SIMPLE FACTORIZATION APPROACH

Small values of P_γ obtained in naive factorization approach have forced theoreticians to search for some enhancement mechanism. The importance of the so-called "nonseparable" contributions given by the diagram of Fig. 2 was stressed in ref.[8,20]. However, as we have already noted, the calculational problems are very complicated here. The main one is to take into account the exchange of the bound quark-antiquark states between the quark or nucleon lines. This problem was treated in ref.[8] with the help of the Bethe-Salpeter equation and the result was enhanced up to $P_\gamma^{WS} = 4.24\times 10^{-8}$ [20] with the RSC potential (the corresponding expression for P_γ was not written explicitly). The standard form of the potential V_W^{PV} can be retained if one changes the coefficients $A, B, \tilde{C}, \tilde{C}', D$ of Eqs. (19,22) connected with those of ref.[8] by Eq. (18) as follows:

$$A \to \cos^2\theta_C \to \frac{1}{2}(1 - 2\sin^2\theta_W),$$

$$B \to 2(1 - 2\sin^2\theta_W) - 2(\cos^2\theta_C - \frac{1}{2} + \sin^2\theta_W),$$

$$\tilde{C} = 0 \to -\frac{2}{\sqrt{3}}\sin^2\theta_W, \qquad (27)$$

$$\tilde{C}' \to -2\sqrt{3}\,\sin^2\theta_W,$$

$$\tilde{D} = 0 \to \frac{1}{2}(\cos^2\theta_C + \frac{1}{2} - \sin^2\theta_W),$$

where we have used Table 2 of ref.[8] with $\zeta_Z = \zeta_W = 1$, Eq. (18) and again put $x^{-1} = 3\sqrt{3}$, $y = \sqrt{3}$ to ensure the absence of the ω-contribution to the separable diagram.

To get insight into the mutual importance of the separable (S) and nonseparable (NS) contributions to P_γ^{WS} and A_γ^{WS}, we substitute the coefficients of Eq. (27) into Eqs. (24-25) and obtain

$$P_\gamma^{HJ} = [(1.6)_S + (1.0)_{NS}] \times 10^{-8} = 2.6 \times 10^{-8},$$

$$A_{pp}^{HJ} = [(-1.80)_S + (4.18)_{NS}] \times 10^{-8} = 2.38 \times 10^{-8}. \qquad (28)$$

One can see that the nonseparable contribution has enhanced the value of P_γ by 60%.

As for P_γ^{tr}, its predicted value is less than that given by Eq. (26) and is equal to 0.96×10^{-6}.

On the other hand Donoghue[9] has shown that the usual factorization approach is insufficient with the Fierz reordering in the quark fields, and that it is necessary to consider the sum of contributions given by the diagrams of Figs. 1,2 to overcome this difficulty. His modified factorization approach has allowed him to overcome the usual difficulties in

calculating the non-separable contributions. The standard form of V_W^{PV} of Eq. (22) can be retained with the changes

$$A \to A - \frac{B}{12} + \frac{\tilde{D}}{36} = \cos^2\theta_C - \frac{1}{6}(1 - 2\sin^2\theta_W),$$

$$B \to B - \frac{2}{3}A + \frac{1}{6}B + \frac{\tilde{\tilde{D}}}{18} = 2(1 - 2\sin^2\theta_W) - \frac{2}{3}(\cos^2\theta_C - \frac{1}{2} + \sin^2\theta_W), \qquad (29)$$

$$\tilde{\tilde{C}} = 0 \to \tilde{C} + \frac{1}{6}(\tilde{C} + \tilde{C}') = -\frac{2}{3\sqrt{3}}\sin^2\theta_W,$$

$$\tilde{C}' \to \tilde{C}' + \frac{1}{6}(\tilde{C} + \tilde{C}') = -\frac{4}{\sqrt{3}}\sin^2\theta_W - \frac{2}{3\sqrt{3}}\sin^2\theta_W,$$

$$\tilde{D} = 0 \to \tilde{\tilde{D}} + 2A + \frac{B}{2} + \frac{1}{6}\tilde{D} = 2(\cos^2\theta_C + \frac{1}{2} - \sin^2\theta_W).$$

We note again that the ω-contribution to P_γ is given entirely by the nonseparable contribution. As a result, Donoghue has obtained $P_\gamma^{RSC} \approx 2.8 \times 10^{-8}$ with Reid-soft-core potential, and $\sin^2\theta_W = 0.35$.

We have tried to separate separable and nonseparable contributions as in the previous case using Donoghue's own expression of P_γ^{RSC}:

$$P_\gamma^{RSC} = (3.84A - 1.5B + 0.1\tilde{D}) \times 10^{-8} =$$

$$= [(1.70)_S + (0.43)_{NS}] \times 10^{-8} = 2.13 \times 10^{-8}, \qquad (30)$$

and we have used $\sin^2\theta_W = 0.25$.

In the same way we have tried to separate S and NS contributions for the asymmetry parameter A_{pp} at $E_{\vec{p}} = 15$ MeV which was calculated in ref.[17] following formulas of ref.[19]

$$A_{pp}(E_{\vec{p}} = 15\ \mathrm{MeV}) = (0.63\cos^2\theta_C - 3.59B -$$

$$- 0.87\tilde{D} - 1.43\tilde{C} - 2.33\tilde{C}') \times 10^{-8} =$$

$$= \begin{cases} [(-1.77)_S + (0.18)_{NS}] \times 10^{-8} = -1.69\times10^{-8} \text{ for } \sin^2\theta_W = 0.25 \\ [(-0.24)_S + (0.50)_{NS}] \times 10^{-8} = 0.26\times10^{-8} \text{ for } \sin^2\theta_W = 0.35 . \end{cases} \quad (31)$$

One can see from Eqs. (30-31) that the enhancement in P_γ does not exceed 30%, while for A_{pp} there are some delicate cancellations. As for P_γ^{tr} with Eq. (29) we obtain 1.2×10^{-6}.

The next step in searching for the enhancement mechanism was the transition to the quantum chromodynamics (QCD). The asymptotically free gluon theory was successfully applied to enhance nonleptonic decay modes in ref. /21/, and this success has inspired the efforts in the analysis of the NN -weak interaction within QCD /22/. The numerical results were obtained in refs. /8,9/. In ref. /8/ the gluon corrections have suppressed the theoretical value of P_γ by a factor of 4, that is $P_\gamma^{QCD} \cong 0.98\times10^{-8}$ instead of $P_\gamma^{WS} = 4.24\times10^{-8}$ without QCD corrections. In the modified factorization approach /9/ the gluon corrections have enhanced the value of P_γ up to $P_\gamma^{QCD} = 6.3\times10^{-8}$ from the value $P_\gamma^{WS} = 2.8\times10^{-}$ Both results are given at $a = 10$, where

$$a = \left[1 + \frac{g}{24\pi^2} (33 - 2n) \ln\left(\frac{M}{\mu}\right) \right] \qquad (32)$$

is the characteristic factor arising from the gluon corrections, g being effective quark-gluon coupling constant, n being the flavour number ($n = 4$ in the GIM model).

One can see that various modifications of the naive factorization approach have not succeeded in increasing the theoretical value of P_γ, which stays at the value $(1\div6)\times10^{-8}$. Moreover, the way of taking into account the nonseparable contributions has proved to be not unique and deserves further investigation. The same is true for the gluon corrections.

4. CONCLUSIONS AND SOME PERSPECTIVES

In the preceding Sections we have outlined the usual factorization approach and its modifications. The main hypothesis of this approach are the validity of saturation of the matrix elements (11) by the vacuum state and the validity of the current-field identity. They are well defined for the naive factorization but their application has no such a solid base while treating nonseparable contributions. At this stage the number of principal and technical assumptions is large, moreover great cancellations occur which make even more difficult to judge whether all sources of enhancement have been already exhausted. Even double-enhancement mechanism, due to nonseparable contributions and due to gluon corrections, does not yield unambiguously large values of P_γ.

These have been the main reasons why we have tried to construct a simple gauge model based on a minimal number of assumptions /10/. Up to now we have succeeded in enhancing essentially the nonseparable contributions based on the vertex diagrams of Fig. 4.

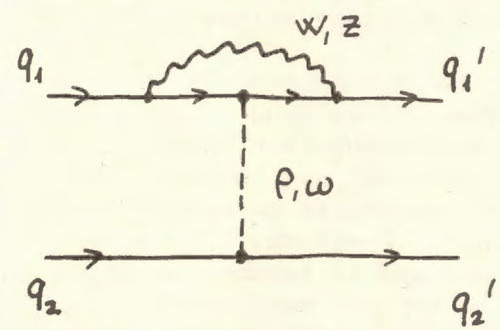

Fig. 4

Updating our results for $\sin^2\theta_W = 0.25$, correcting wrong sign in front of A, and changing the definitions by more suitable ones to compare with the usual nonseparable contributions, we wrote Eq. (9) of ref. /10 in the form

$$A = \frac{1}{2}(1 - 2\sin^2\theta_W), \quad B = 2\left(\cos^2\theta_C - \frac{1}{2} + \sin^2\theta_W\right), \quad \tilde{C} = \frac{2}{\sqrt{3}}\sin^2\theta_W,$$

$$\tilde{C}' = \frac{4}{\sqrt{3}} \sin^2\theta_W \frac{g_{\omega NN}}{g_{\rho NN}}, \quad \tilde{D} = -\frac{g_{\omega NN}}{g_{\rho NN}} (\cos^2\theta_C + \frac{1}{2} - \sin^2\theta_W), \tag{33}$$

$$f_V = -\sqrt{2}\, g_{\rho NN}\, a_W /10/ = -3.42\times10^{-6} \quad * .$$

In /10/ $\xi = 1$ was implicitly assumed as we could use usual formulas for ξ within our quark hypothesis. But it would be more reasonable to use it as free parameter say in the range from $(6\sqrt{3})^{-1}$ to 1. With Eq. (33) the nonseparable contributions for P_γ now yield $P_\gamma^{HJ} = -5.52\times10^{-8}$ for Hamada-Johnston's potential, $P_\gamma^{RSC} = 3.22\times10^{-8}$ for Reid's soft-core one, $P_\gamma^{T} = 3.29\times10^{-8}$ for Tamagaki's one. One can see that our nonseparable contribution is significantly enhanced. Moreover, naive sum with the usual separable contribution gives $P_\gamma \sim 10^{-7}$.

Substituting Eq. (33) into Eq. (25), we obtain the nonseparable part of $A_{pp}^{HJ} = -1.16\times10^{-8}$.

We consider our results quoted here as preliminary, as they are based mainly on the expressions from /12/ which are very convenient to use but where we have met some unclear points concerning the D-wave contribution. Indeed, in /23/ it was shown that different input of the D-wave could give deviations of P_γ by an order of magnitude.

Nevertheless, the situation encourages us to try to construct the field-theoretical gauge model which would describe the weak nucleon-nucleon interactions in a more rigorous and consistent way. However, we hardly expect that it is possible to exceed the value $P_\gamma \sim 10^{-7}$ along these lines. Therefore, if the result $P_\gamma^{exp} \sim 10^{-6}$ would be confirmed in future, we would arrive at the necessity to explore such exotic hypothesis as say contact weak interaction in the compact 6-quark states /24/ or its inverse alternative-long-range weak interaction through two-pion exchanges /25/. In all cases it would be useful also to continue phenomenological calculations of P_γ along the lines of /15/ within the scheme of relativistic deuteron.

* The values of A, B, $\tilde{C}$, $\tilde{C}'$, $\tilde{D}$ for the various models are given in Table 1.

Table 1

W–S	f_V	A	B	$\tilde{C}$	$\tilde{C}'$	$\tilde{D}$
S /12/	-2.127×10^{-6}	c^2	$2(1-2s^2)$	0	$-\frac{4}{\sqrt{3}} s^2$	0
S+NS /8/	-2.127×10^{-6}	$c^2 - \frac{1}{2}(1-2s^2)$	$2(1-2s^2) - 2(c^2 - \frac{1}{2} + s^2)$	$-\frac{2}{\sqrt{3}} s^2$	$-2\sqrt{3}\, s^2$	$\frac{1}{2}(c^2 + \frac{1}{2} - s^2)$
S+NS /9/	-2.127×10^{-6}	$c^2 - \frac{1}{6}(1-2s^2)$	$2(1-2s^2) - \frac{2}{3}(c^2 - \frac{1}{2} + s^2)$	$-\frac{2}{3\sqrt{3}} s^2$	$-\frac{14}{3\sqrt{3}} s^2$	$2(c^2 + \frac{1}{2} - s^2)$
NS /10/	-3.42×10^{-6}	$\frac{1}{\sqrt{2}}(1-2s^2)$	$2(c^2 - \frac{1}{2} + s^2)$	$\frac{2}{\sqrt{3}} s^2$	$\frac{4}{3}\frac{g_{\omega NN}}{g_{\rho NN}} s^2$	$-\frac{g_{\omega NN}}{g_{\rho NN}}(c^2 + \frac{1}{2} - s^2)$

$c \equiv \cos\theta_C$

$s \equiv \sin\theta_W$

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