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A b s t r a c t

The radiative corrections to a parity violating electron--quark interaction are calculated in the Weinberg-Salam model. The strong interactions are taken into account in the frame of Quantum Chromodynamics.

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1. Introduction

In the Weinberg-Salam theory the parity violating electron-quark interaction is due to the Z-boson exchange (Fig. 1) and is described by two spatial structures: the product of the axial electron current by the hadron vector one and the product of the vector electron current by the hadron axial on. One can easily verify that at small transfer ($K^2 \ll M_Z^2$) the corresponding effective Hamiltonian has the form:

$$H = -\frac{G}{\sqrt{2}}\left[\varkappa_1 \bar{e}\gamma_\mu\gamma_5 e\,\bar{q}\gamma_\mu q + \varkappa_2 \bar{e}\gamma_\mu e\,\bar{q}\gamma_\mu\gamma_5 q\right]$$

$$\varkappa_1 = T_3 - 2Q\sin^2\theta_W, \quad \varkappa_2 = T_3(1-4\sin^2\theta_W) \qquad (1)$$

where Q and T_3 are the charge and projection of a quark isospin, θ_W is the Weinberg angle - a free parameter of the model. The interaction associated with the first structure was observed in an atomic experiment /1/ and in the deep-inelastic electron-proton and electron-deuteron scattering /2/. The experimental value of $\varkappa_1$ agrees with the theoretical one. To measure the constant $\varkappa_2$ is a more complicated problem even if one assumes that $\varkappa_2 \sim 1$. Meanwhile $\varkappa_2$ is strongly suppressed in the theory since experimental value of $\sin^2\theta_W$ is close to 1/4 ($\sin^2\theta_W = 0.230 \pm 0.015$ /3/). From the one hand, this circumstance hinders the observation of the relevant effects but, from the other hand, a new possibility arises: the precise measuring of $\sin^2\theta_W$. In this situation the question concerning the calculations of corrections to $\varkappa_2$ acquires especial interest.

The isoscalar part of the correction $\delta\varkappa_2$ due to the second order to an electro-weak interaction was earlier calculated in Ref. /4/ (see also /5/). The influence of strong interactions on this correction was not discussed. In the present paper the isoscalar and isovector parts of the correction $\delta\varkappa_2$ due to the second order to electro-weak interaction are calculated. The influence of strong interactions is taken into account. The quantity δ_{eff} is also found.

2. One-loop corrections

The Lagrangian of electro-weak interactions in the Weinberg-Salam model has the form:

$$\mathcal{L} = e A_\mu J_\mu^{em} + \frac{g}{\sqrt{2}}\left(W_\mu J_\mu^{+} + h.c.\right) + \frac{g}{\cos\theta_W} Z_\mu J_\mu^{0}$$

$$J_\mu^{em} = \bar e \gamma_\mu e - \bar q Q \gamma_\mu q$$

$$J_\mu^{+} = \bar q \gamma_\mu (1+\gamma_5) T^{+} q + \bar\nu \gamma_\mu (1+\gamma_5) e \qquad (2)$$

$$J_\mu^{0} = \bar q \gamma_\mu \left[T_3 (1+\gamma_5) - 2Q\sin^2\theta_W\right] q +$$

$$+ \bar e \gamma_\mu \left[\left(-\frac{1}{2} + 2\sin^2\theta_W\right) - \frac{1}{2}\gamma_5\right] e$$

The matrix of quark charges Q and matrices T_3 and T^+ look as follows:

$$Q = \begin{pmatrix} \frac{2}{3} & & & \\ & \frac{2}{3} & & \\ & & -\frac{1}{3} & \\ & & & -\frac{1}{3} \end{pmatrix}, \quad T_3 = \frac{1}{2}\begin{pmatrix} 1 & & & \\ & 1 & & \\ & & -1 & \\ & & & -1 \end{pmatrix}, \quad T^{+} = \begin{vmatrix} 0 & 0 & -\sin\theta_c & \cos\theta_c \\ 0 & 0 & \cos\theta_c & \sin\theta_c \\ 0 & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 \end{vmatrix}$$

$$T^{+} = T_1 + iT_2, \quad T^{-} = T_1 - iT_2$$

where θ_c is the Cabibbo angle. We consider four quarks only: $q = (c, u, d, s)$. As for heavy quarks, they influence the corrections under discussion very slightly. We carry out our calculations within logarithmic accuracy, i.e. the corrections $\sim \alpha \ln \frac{M_W^2}{k^2}$ are calculated. The graphs giving a logarithmic contribution are drawn in Figs. 2a,b,c. The diagram with the double W-boson exchange (Fig. 2d) should be taken into account also: it does not give a logarithm but has the enhancement factor $1/\sin^2\theta_W$.

Let us discuss the diagram in Fig. 2c. It contains the logarithmically divergent loop (Fig. 3) which should be normalized at some point. As Fig. 3 corresponds to the γ-Z- mixing, it is clear that it is the renormalization of $\sin^2\theta_W$. In principle, $\sin^2\theta_W$ may be normalized at any point, for example at $k^2 = M_Z^2$. But at present $\sin^2\theta_W$ is determined from the experiments on neutrino-hadron and electron-hadron scattering at comparatively small transfers $k^2 \sim 1\,\text{GeV}^2$. In view of this, the normalization of $\sin^2\theta_W$ at $k^2 = 0$ seems quite natural. Thus, we determine a renormalized quantity of $\sin^2\theta_W$ so that the amplitude associated with Fig. 3 is zero at $k^2 = 0$. At such a definition, in any process where $\sin^2\theta_W$ is measured, for example in the deep-inelastic νN scattering, it is necessary to take into account, besides the tree graph, both the correction related to Fig. 3 and the radiative corrections to vertex. However, all these corrections now are smaller than the experimental errors.

Let us return to graph 2c. With the said above, its contribution is zero at $k^2 = 0$ while at large k^2 grows as $\sim \ln \frac{k^2}{\mu_0^2}$. For the quark loop $\mu_0 \sim m_\pi + m_\rho$ is an infrared parameter. For the lepton loop μ_0 should be replaced by the lepton mass m_e. But the contribution of the lepton loops is proportional to $(1-4\sin^2\theta_W)$ and therefore it can be neglected.

Calculation of the down block of diagrams 2b,c which corresponds to the quark anapole moment gives rise, within logarithmic accuracy, to the following result:

$$\frac{e}{16\pi^2}\frac{G}{\sqrt{2}}\left(k^2 g_{\mu\nu} - k_\mu k_\nu\right)\left[\bar q \gamma_\nu \gamma_5 \left(C_S(k^2) + C_V(k^2) T_3\right) q\right]$$

$$C_S(k^2) = \left(\frac{1}{3} - \frac{4}{9}\sin^2\theta_W\right)\ln\frac{M_W^2}{(k^2, \mu_0^2)}$$

$$C_V(k^2) = \left(-\frac{2}{3} - \frac{40}{27}\sin^2\theta_W\right)\ln\frac{M_W^2}{(k^2, \mu_0^2)} - \qquad (3)$$

$$-8\left(1 - \frac{20}{9}\sin^2\theta_W\right)\ln\frac{(k^2, \mu_0^2)}{\mu_0^2}$$

$$(k^2, \mu_0^2) \equiv \max(k^2, \mu_0^2)$$

The correction to the constant $\varkappa_2$ has the form:

$$\delta\varkappa_2 = \delta\varkappa_{2,s} + \delta\varkappa_{2,v} T_3$$

$$\delta\varkappa_{2,s} = -\frac{\alpha}{4\pi}\Big[\Big(-\frac{3}{2} + \frac{10}{3}\sin^2\theta_w\Big)\ln\frac{M_w^2}{k^2} - \frac{3}{4\sin^2\theta_w} + C_s(k^2)\Big] \qquad (4)$$

$$\delta\varkappa_{2,v} = -\frac{\alpha}{4\pi}\Big[\big(-1 + 4\sin^2\theta_w\big)\ln\frac{M_w^2}{k^2} - \frac{5}{2\sin^2\theta_w} + C_v(k^2)\Big]$$

where $\varkappa_s$ and $\varkappa_v$ are isoscalar and isovector parts of the $\varkappa$. It is worth noting that the formulae (3) and (4) at $k^2 < m_c^2$ are valid only for light u, d, s - quarks.

The isoscalar part of the correction (4) at $k^2 = 0$ coincides with that derived in Ref. /4/.

Unlike $\varkappa_2$, the constant $\varkappa_1$ in the Hamiltonian (1) is not small. But for completeness we calculate also the correction $\delta\varkappa_1$:

$$\delta\varkappa_1 = \delta\varkappa_{1,s} + \delta\varkappa_{1,v} T_3$$

$$\delta\varkappa_{1,s} = -\frac{\alpha}{4\pi}\Big[-\frac{3}{4\sin^2\theta_w} - \frac{2}{3}\Big(1 - \frac{20}{9}\sin^2\theta_w\Big)\ln\frac{(k^2, \mu_0^2)}{\mu_0^2}\Big] \qquad (5)$$

$$\delta\varkappa_{1,v} = -\frac{\alpha}{4\pi}\Big[-\frac{5}{2\sin^2\theta_w} - 4\Big(1 - \frac{20}{9}\sin^2\theta_w\Big)\ln\frac{(k^2, \mu_0^2)}{\mu_0^2}\Big]$$

This correction is due to graphs 2d and 4. The contributions of the rest graphs are proportional to $1 - 4\sin^2\theta_w$ and they are omitted.

3. The influence of strong interactions

The strong interactions are taken into account in the frame of Quantum Chromodynamics. This is done, as usually, within logarithmic accuracy, i.e. we consider corrections $\sim \alpha_s \ln\frac{M_w^2}{k^2}$, $(\alpha_s \ln\frac{M_w^2}{k^2})^2$ and so on. It is most convenient to work in the Landau gauge over gluons. One can easily see that the gluon corrections in the diagrams in Figs. 2a, d as well as in Fig. 4 do not give the logarithms, i.e. the strong interactions do not affect the relevant amplitudes within accepted accuracy. Meanwhile in the diagrams Fig. 2b, the gluons give the logarithmically large corrections $\sim \alpha_s \ln\frac{M_w^2}{k^2}$. Hence, strong interactions affect the magnitude of $\delta\varkappa_2$ and do not affect $\delta\varkappa_1$.

Let us cut the photon line in the diagrams 2b,c and calculate corrections to the down block that corresponds to the anapole moment of the quark. As the parameter $\alpha_s \ln\frac{M_w^2}{k^2} \sim 1$, summation of all the orders of perturbation theory is necessary. To do this, we'll use the renormalization group technique.

We consider the case of comparatively small transfer $k^2 \ll M_w^2$. In this situation it is convenient to work with the effective four-fermion interaction. The divergent integrals should be cut, of course, on M_w. Besides the operators which enter the bare effective Hamiltonian, some new operators arise due to radiative corrections. We are interested in the processes without change of strangeness and charm ($\Delta S = \Delta C = 0$). Taking into account this circumstance let us list independent operators of minimal dimension which can be involved in the violating parity effective Hamiltonian.

Operators that correspond to the anapole moment of quarks have the dimension 6 and look as follows:

$$O_s = -e\,\partial_\mu F_{\mu\nu}\,\bar q \gamma_\nu \gamma_5 q\,, \quad O_v = -e\,\partial_\mu F_{\mu\nu}\,\bar q T_3 \gamma_\nu \gamma_5 q \qquad (6)$$

$F_{\mu\nu}$ - is the electromagnetic field strength tensor. We include the electron charge in the definition of operators since they always enter in the effective Hamiltonian multiplied by it. The other pseudoscalar gauge invariant operators of dimention 6 which enter the effective Hamiltonian, has the form:

$$O(M,\Lambda) = \bar{q}\gamma_\mu\gamma_5 M q\,\bar{q}\gamma_\mu \Lambda q \tag{7}$$

where M and Λ are arbitrary matrices acting in $SU(4)$ and $SU(3)^c$ spaces. The remaining operators of the same dimension can be reduced, by means of the equations of motion, to the operators listed above. For example, the operator

$$\mathcal{D}_\nu G^a_{\mu\nu}\,\bar{q}\gamma_\mu\gamma_5\lambda^a q \tag{8}$$

where λ^a are the $SU(3)^c$ matrices, $Sp\,\lambda^a\lambda^b = 2\delta^{ab}$, $G^a_{\mu\nu}$ is the gluon field strength tensor. It can be reduced to the operators (7). The single operator of dimension 5 is as follows:

$$G^a_{\mu\nu}\,\bar{q}\,\sigma_{\mu\nu}\gamma_5\lambda^a q \tag{9}$$

However, it is forbidden due to CP-invariance.

In the renormalization of the four-quark operators (7), we follows Ref. /6/. According to /6/, these operators can be classified over the $SU(4)$ - group representations. Operators belonging to the 84 - and 20 - dimensional $SU(4)$ representations have the following form:

$$O_{84}(T,T) = \frac{2}{3}\left[O(T,T) - \frac{1}{20}O(1,1)\right] + \frac{1}{4}\left[O(T\lambda^a, T\lambda^a) - \frac{1}{20}O(\lambda^a,\lambda^a)\right],$$

$$O_{20}(T,T) = \frac{1}{3}\left[O(T,T) + \frac{1}{12}O(1,1)\right] - \frac{1}{4}\left[O(T\lambda^a, T\lambda^a) + \frac{1}{12}O(\lambda^a,\lambda^a)\right], \quad T = T_1, T_2, T_3 \tag{10}$$

The index denotes the $SU(4)$ - group representation to which the operator belongs. Matrices T_i are introduced in formula (2). For the four-quark interaction, the following operators appear:

$$O_{1,2} = O_{84,20}(T_1,T_1) + O_{84,20}(T_2,T_2) + (1-2\sin^2\theta_W)\,O_{84,20}(T_3,T_3)$$

$$O_3 = O(1,1)$$

$$O_4 = O(\lambda^a,\lambda^a)$$

$$O_5 = O(T_3, 1), \qquad O_7 = O(1, T_3)$$

$$O_6 = O(T_3\lambda^a, \lambda^a), \qquad O_8 = O(\lambda^a, T_3\lambda^a) \tag{11}$$

Operators O_3 and O_4 are $SU(4)$ singlets, while $O_5 \div O_8$ correspond to 15 dimensional representation.

The violating parity effective Hamiltonian can be expanded in terms of operators O_5 O_6, $O_1 \div O_4$:

$$H = \frac{G}{\sqrt{2}}\sum c_i O_i$$

The coefficients in this expansion depend on the external momenta μ^2. At $\mu^2 \sim M_W^2$ the coefficients c_i are not renormalized by strong interactions and are determined from the comparison with the bare Hamiltonian:

$$C(M_W) = \left(0,0,2,2,\ \frac{3-2\sin^2\theta_W}{90},\ \frac{3-2\sin^2\theta_W}{15},\ -\frac{2}{3}\sin^2\theta_W,\ 0,0,0\right) \tag{12}$$

We are interested in the Hamiltonian at $\mu^2 \ll M_W^2$. In this region the coefficients C_i and the constant of strong interactions α_s satisfy the renormalization-group equations (see, for example Ref. /7/):

$$\mu \frac{d\alpha_s(\mu)}{d\mu} = -\frac{b}{2\pi}\alpha_s^2(\mu)$$

$$\mu \frac{dC_i(\mu)}{d\mu} = -\frac{\alpha_s}{2\pi}\gamma_{ij} C_j(\mu) \tag{13}$$

where $b = 11 - \frac{2}{3}N$, N is the number of flavours In calculating the effective Hamiltonian the regions $M_W \div m_c$ and $m_c \div \mu_0$ should be distinguished (m_c is the mass of c - - quark). Strictly speaking, our consideration concerns the first region where we can neglect the mass of the c-quark. Meanwhile in the region $\mu < m_c$ three light quarks should be taken into consideration only. However, the absence of the c-quark influences slightly the coefficients for the operators which contain the fields of light quarks only. An accurate calculation with the intermediate scale m_c shows that the coefficient C_V differs not more 3% from that calculated without this scale. As for C_S , the difference is about 50%. Such large difference is connected with the strong compensation of different terms. In any case C_S is small, and it is therefore calculated unreliably. In view of this, we do not take into account, for simplicity, the intermediate scale m_c, We assume that the $SU(4)$ symmetry is true up to μ_0. The electromagnetic interaction is taken into account in the first order only. Therefore O_S and O_V do not affect the renormalization of operators $O_1 \div O_8$. It is easy to see that the anomalous dimensions of operators O_S and O_V are zero. The mixing coefficients of O_S, O_V with operators $O_1 \div O_8$ are determined by the diagram in Fig. 5 and are as follows:

$$\gamma_{Si} = \left(0, 0, \frac{26}{9}, \frac{32}{27}, \frac{1}{3}, \frac{16}{9}, \frac{13}{3}, \frac{16}{9}\right) \tag{14}$$

$$\gamma_{Vi} = \left(\frac{8}{5}(1-4\sin^2\theta_W), \frac{4}{3}(1-\frac{4}{3}\sin^2\theta_W), \frac{4}{3}, \frac{64}{9}, \frac{26}{9}, \frac{32}{27}, \frac{2}{9}, \frac{32}{27}\right)$$

$$i = 1 \div 8$$

The matrix of anomalous dimensions of four-quark operators $O_1 \div O_8$ is determined by the graphs in Fig. 6 and it can be represented as a series of matrices corresponding to different representation- of the $SU(4)$ group /6/.

$$O_1: \ \gamma = -2, \quad O_2: \ \gamma = 4$$

$$O_3, O_4: \quad \gamma = \begin{pmatrix} 0 & -\frac{16}{3} \\ -\frac{11}{6} & -\frac{4}{9} \end{pmatrix}$$

$$O_5 \div O_8: \quad \gamma = \begin{pmatrix} 0 & 0 & 0 & -\frac{16}{3} \\ -\frac{1}{3} & \frac{37}{18} & -\frac{11}{6} & -\frac{41}{18} \\ 0 & -\frac{16}{3} & 0 & 0 \\ -\frac{3}{2} & -\frac{5}{2} & 0 & \frac{9}{2} \end{pmatrix} \tag{15}$$

The solution of eqs. (13) is of the form:

$$C(t) = \left(\frac{\alpha_s(t)}{\alpha_s(0)}\right)^{\frac{\gamma}{b}} C(t=0) \qquad i = 1 \div 8$$

$$C_{S,V} = \int_0^t \gamma_{S,V;i}\, C_i(t')\, dt' \tag{16}$$

$$\alpha_s(t) = \frac{\alpha_s(0)}{1 - \frac{b}{4\pi}\alpha_s(0)\, t}, \quad \alpha_s(0) = \frac{1}{1 + \frac{b}{4\pi}\ln\frac{M_W^2}{\mu_0^2}}$$

where $t = \ln\frac{M_W^2}{\mu^2}$.

In order to obtain a correct answer for a quark anapole moment it remains to normalize $\sin^2\theta_W$. We keep in mind the following circumstance. Our calculations with effective Hamiltonian corresponds to the normalization of $\sin^2\theta_W$ at $K^2=M_W^2$. To proceed to the accepted normalization, it is necessary to subtract from the $C_V(t)$ the contribution of the graph Fig. 2b (without the electron line) coming from the region $\mu_0^2 \div M_W^2$.

$$C_V(K^2)\to C_V(K^2)-8\left(1-\frac{20}{9}\sin^2\theta_W\right)\ln\frac{M_W^2}{\mu_0^2} \tag{17}$$

In this formula strong compensations take place. For example, at $K^2=0$, $C_V(0)\to -\frac{1}{20}C_V(0)$. (We don't give here analytical formulae for C_V and C_S taking into account strong interactions because of their cumbersome character). It is worth stressing that such a compensation does not implies that the quantity $C_V(0)$ is calculated with bad accuracy. The point is that the substraction corresponds to the exact vanishing of the contribution of graph 2c. In view of this, the accuracy of calculation of $C_V(0)$ is determined by the accuracy of calculation for all the remaining graphs and the error can be expected to be not more than 10-30%. At $K^2\sim$ 1÷10 GeV the accuracy of calculation of $C_V(K^2)$ is worse. This is due to the fact that in this region the main contribution to $C_V(K^2)$ comes from graph 2c which is calculated within logarithmic accuracy $\sim\ln\frac{K^2}{\mu_0^2}$. Since the logarithm is not yet large here, the error can be of the order of 50%. If one subtracts from $C_V(K^2)$ the contribution of diagram 2c

$$\tilde C_V(K^2)=C_V(K^2)+8\left(1-\frac{20}{9}\sin^2\theta_W\right)\ln\frac{(K^2+\mu_0^2)}{\mu_0^2} \tag{18}$$

the difference $\tilde C(K^2)$ for all $K^2\ll M_W^2$ is calculated with an accuracy of about 10-30%. The graph $\tilde C_V(K^2)$ is given in Fig. 7. As for the quantity $C_S(K^2)$, it is small, as mentioned above, because of strong compensation. Therefore the accuracy of calculation is $\sim$ 50% at $K^2=0$ and somewhat higher at $K^2\sim$ 1-10 GeV. The graph $C_S(K^2)$ is given in Fig. 8. It is easy to find the correction $\delta\mathcal{X}_2$, substituting $C_S(K^2)$ and $C_V(K^2)$ in the formula (4). Note, that although the accuracy of cal-

culations of $C_S(K^2)$ is $\sim$ 50%, the quantity $\delta\mathcal{X}_{2,S}$ is calculated with an accuracy of 10-30%, because the contribution of $C_S(K^2)$ to $\delta\mathcal{X}_{2,S}$ is small.

4. Conclusion

At $\sin^2\theta_W=0.236$ and $K^2=0$ the constant $\mathcal{X}_2$ and the corrections found by us are as follows:

$$\mathcal{X}_{2,S}=0 \qquad \mathcal{X}_{2,V}=0.68$$

$$\delta\mathcal{X}_{2,S}=0.6\cdot10^{-2},\quad \delta\mathcal{X}_{2,V}=0.9\cdot10^{-2}$$

with taking into account strong interactions

$$\delta\mathcal{X}_{2,S}=0.5\cdot10^{-2},\quad \delta\mathcal{X}_{2,V}=1.3\cdot10^{-2} \tag{19}$$

without taking into account strong interactions.

Besides the corrections calculated by us $\sim\alpha$, there also exists the correction $\sim\alpha_s^2$ connected with graph 9. Its calculation has been carried out in Ref. /8/. The result is as follows:

$$\delta\tilde{\mathcal{Z}}_{2,S}\simeq(1-4\sin^2\theta_W)\frac{1}{4}\,\frac{\alpha_s(m_c^2+K^2)}{\pi}\,\frac{\alpha_s(m_s^2+K^2)}{\pi}\ln\frac{m_c^2+K^2}{m_s^2+K^2} \tag{20}$$

This correction does not contain α but it is proportional to ($1-4\sin^2\theta_W$). At $\sin^2\theta_W=0.23$ and $K^2=0$, $\delta\tilde{\mathcal{Z}}_{2,S}=0.2\cdot10^{-2}$ it is less approximately by a factor of 3 compared to the contribution calculated in this work. Moreover, the correction (20) decreases fastly with K^2.

In a real experiment the electron interacts with the nucleon or nucleus. The question then poses in what cases the calculated correction may be significant, and when it is much more less than the corrections associated with the nucleon or nucleus structure. It is clear that for the deep-inelastic scattering there are no other corrections, besides the found one, since all the structural effects are suppressed by the powers μ_0^2/K^2.

In the elastic electron-nucleon scattering and hence in atomic experiments there is one more mechanism leading to parity violation (Fig. 10). It is obvious that one can compare only the matrix elements with respect to nucleon from the effective Hamiltonian (1) and diagram 10. These matrix elements are connected with large distances. When there is one nucleon in an intermediate state, the contribution from diagram 10 to the isoscalar part $\mathcal{F}_Z$, calculated in the bag model, is one order less than the contribution associated with short distances. As shown in Ref. /9/, the main contribution is connected with $N\pi$ intermediate states and it is comparable with that we have calculated.

At last, in atomic experiments there is the contribution arising due to the interaction of nucleus nuclons /10/. Apparently, the main correction in the experiments with heavy atoms or molecules is caused by this mechanism. For deuteron this correction is comparable with that considered in the present paper.

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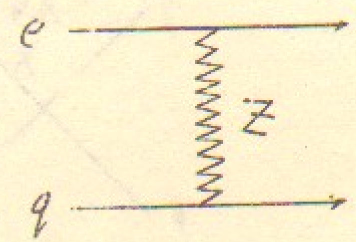
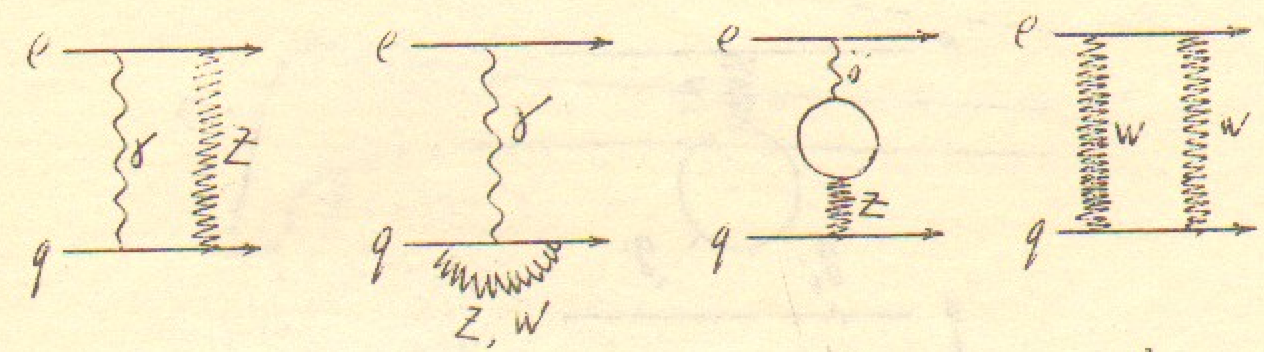


Fig. 1



a b c d

Fig. 2

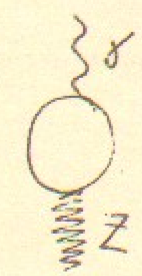


Fig. 3

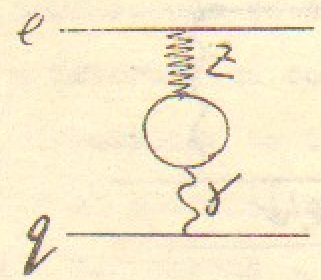


Fig. 4

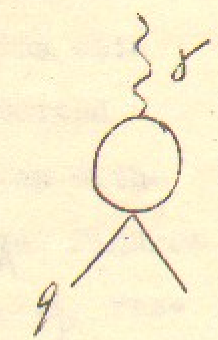


Fig. 5

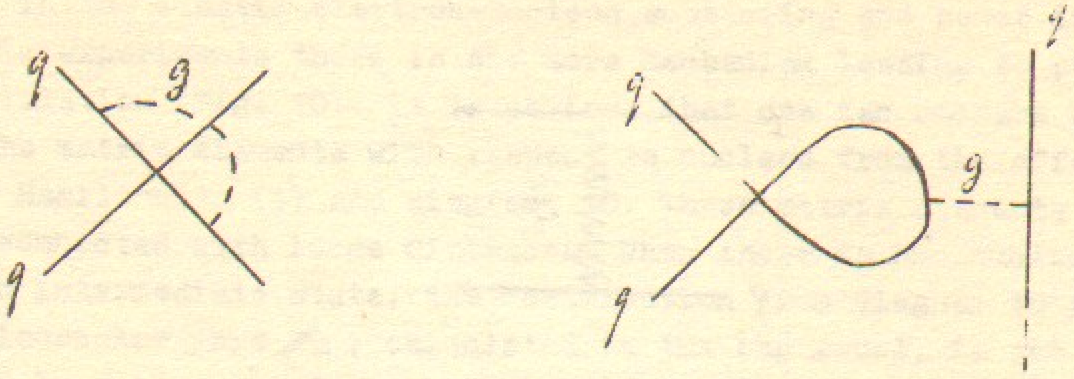


Fig. 6

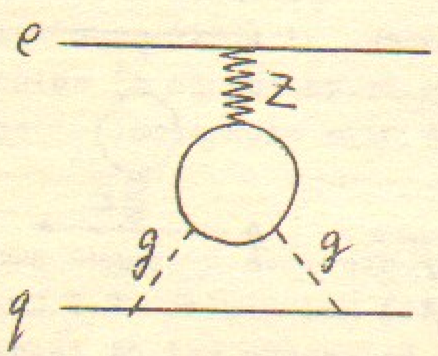


Fig. 9

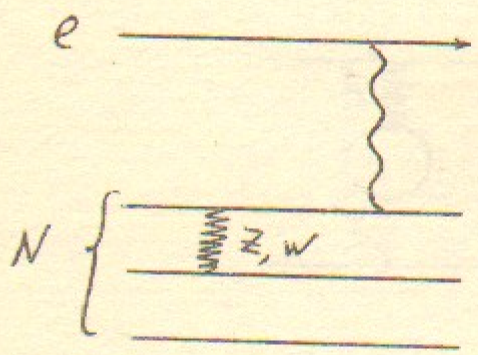


Fig. 10

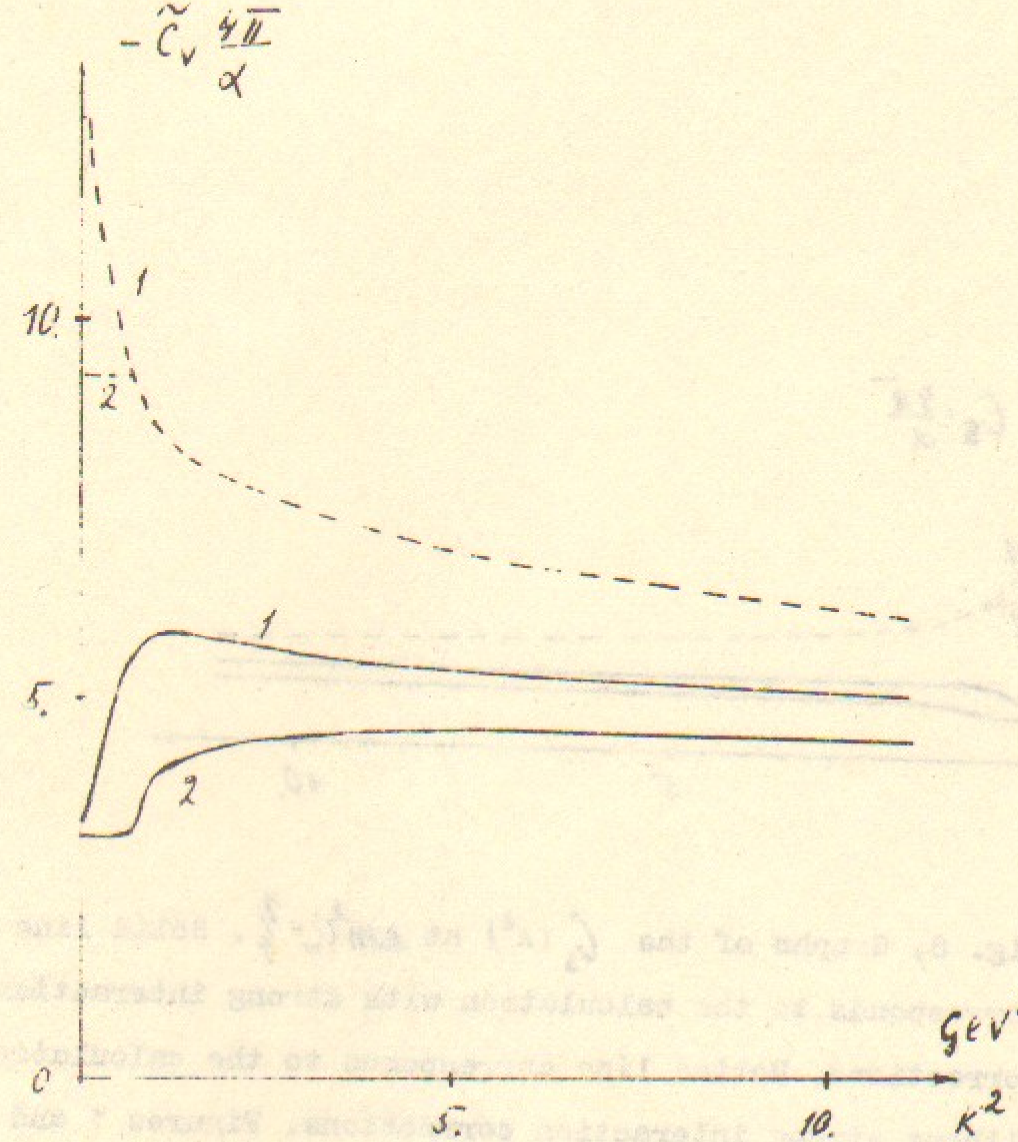


Fig. 7. Graphs of the $\tilde{C}_V(\kappa^2)$ at $\sin^2\theta_W = 1/4$. Solid line corresponds to the calculation with strong interaction corrections. Dotted line corresponds to the calculation without strong interaction corrections. Figures 1 and 2 correspond $\mu_0 = m_\pi$ and $\mu_0 = m_\rho$ respectively.

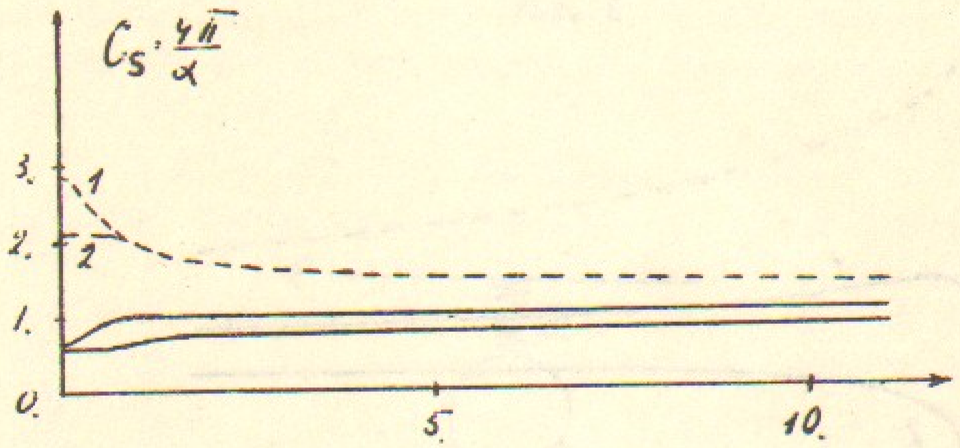


Fig. 8. Graphs of the $C_s(k^2)$ at $\sin^2\theta_W = \frac{1}{4}$. Solid line corresponds to the calculation with strong interaction corrections. Dotted line corresponds to the calculation without strong interaction corrections. Figures 1 and 2 correspond $\mu_0 = m_\pi$ and $\mu_0 = m_\rho$ respectively.