Partial conservation of the axial current and nuclear screening of the interaction of high-energy neutrinos

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The consequences of the PCAC hypothesis for high energy physics are considered. It is shown that, in contrast to the generally accepted point of view, the space-time pictures of nuclear screening are different in the cases of the neutrino and the photon. The main reason for the screening is not the virtual emission of long-lived hadron fluctuations by the high-energy neutrino, but the real process of diffractive neutrino production of pions on nuclear nucleons. This results in saturation of nuclear screening already at low energies.

1. HADRONIC PROPERTIES OF LEPTONS AND PHOTONS

A remarkable event has recently occurred and has, regrettably, remained unnoticed: the WA59 collaboration has, by combining the statistics obtained in the ν and $\bar{\nu}$ beams in the BEBC chamber, confirmed for the first time[1,2] the screening of the neutrino interaction with nuclei at low values of Q^2. What is first of all extraordinary about this result is that the phenomenon of nuclear screening of the interaction involving the weak axial current was predicted by Bell[3] 25 years ago. It was precisely these ideas that initiated the development of the vector-dominance model and the study of photo-nuclear reactions.[4,5] Thereafter nuclear screening of the vector current was reliably confirmed in experiments with real, as well as virtual, photons (see, e.g., the review Ref. 6). However, the problem of experimental observation of the screening of the axial current turned out to be so complicated that its solution required a quarter of a century.

The reasons that prevent particles with small interaction cross sections (photon, leptons) from penetrating the nucleus deeply, are quite interesting. The following space-time picture is usually associated with the process: the lepton (Fig. 1) or photon (Fig. 2) may virtually emit a hadronic fluctuation for a time interval related by the uncertainty relation and Lorentz transformations to the energy ν of the fluctuation, its mass M and the square of its 4-momentum Q^2:

$$t\approx\frac{2\nu}{M^2+Q^2}. \qquad (1)$$

Although the hadronic fluctuation is present in the lepton (photon) wave function with a small weight, it may live a long time and interact with a group of nucleons. Since the interaction cross section of such a rare hadronic fluctuation is large, the nucleons should screen each other. It is widely asserted that the nuclear screening of leptons (photons) is saturated at energies for which the lifetime (1) of the fluctuation substantially exceeds the dimensions of the nucleus. In fact the saturation condition is different: $t\gg\lambda$, where $\lambda=2/\sigma\rho$ is the length of the hadron mean free path in the nucleus, σ is the total hadron-nucleon interaction cross section, and ρ is the nuclear density. Indeed, the given nucleon is completely screened by a layer of nuclear matter of thickness of order λ, so that even if the lifetime of the fluctuation is less than the nuclear radius, but $t\gg\lambda$ holds, an increase in the energy ν will not strengthen the screening.

2. PCAC AND THE NONTRIVIAL PICTURE OF AXIAL CURRENT SCREENING

In this section we show that the space-time picture of the interaction described above is not applicable to the description of nuclear screening of the weak axial current.[7] This assertion appears paradoxical since it is based precisely on the nuclear screening of the axial current predicted by Bell. However, a more detailed study of this question also results in paradoxes.

Bell started in his work from the Adler relation,[8] which connects the cross section for the process $\nu T\to lF$ at $Q^2=0$ with the cross section for the process $\pi T\to F$, where T and F are the target and the final hadronic state, and $-Q^2$ is the square of the 4-momentum transferred from the neutrino ν to the lepton l:

$$\left.\frac{d^2\sigma(\nu T\to lF)}{dQ^2\,d\nu}\right|_{Q^2=0}=\frac{G^2}{2\pi^2}f_\pi^2\left[\frac{1}{\nu}-\frac{1}{E}\right]\sigma(\pi T\to F). \qquad (2)$$

Here E is the neutrino energy, the difference between the energies of the neutrino and the lepton is denoted as usual by the letter ν, G is the Fermi constant of weak interactions, and f_π is the decay constant for $\pi\to\mu\nu$. Here we omit terms $\sim o(m_l^2)$, where m_l is the mass of the lepton.

This relation is a direct consequence of the PCAC hypothesis, according to which the divergence of the hadronic axial current is proportional to the pion field:

$$q_\mu A_\mu=f_\pi m_\pi^2\Phi_\pi.$$

It is precisely the smallness of the pion mass which permits one to speak of approximate conservation of the current. In those cases where it does not matter we shall pass to the chiral limit $m_\pi=0$, i.e., view the axial current as exactly conserved (of course to within the axial anomaly), in order to simplify the exposition.

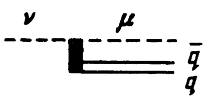

FIG. 1. Virtual quark-antiquark pair production by a neutrino.

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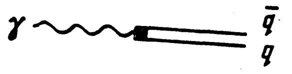

FIG. 2. Virtual transition of a photon into a quark-antiquark pair.

A natural, but incorrect, association with the Adler relation (2) is the dominance of the one-pion exchange diagram shown in Fig. 3. That this interpretation is wrong can be seen by writing the contribution of this diagram to the amplitude for the process in the form

$$A(\nu T\to lF)=\frac{G}{2^{1/2}}f_\pi L_\mu\frac{q_\mu}{Q^2+m_\pi^2}A(\pi T\to F). \tag{3}$$

Here $L_\mu=\bar{l}\gamma_\mu(1+\gamma_5)\nu$ is the lepton current, which is transverse to within the lepton mass: $L_\mu q_\mu=0$. It follows that the contribution (3) is absent, i.e., the neutrino cannot emit a virtual pion in vacuum.[9,10]

Thus the process $\nu\to lF$ is determined by heavier hadronic fluctuations, for example, a_1 mesons. It should then be expected that the saturation of nuclear screening would occur at energies $\nu\gg\lambda m_a^2/2\approx 20$ GeV. At such energies the total cross section should depend on A like $A^{2/3}$. On the other hand, if the lifetime of the fluctuation is less than the mean inter-nucleon distance in the nucleus, $d\approx 2$ fm, i.e., $\nu\ll dm_a^2/2\simeq 6$ GeV, the screening disappears and $\sigma_{\rm tot}(\nu A)\propto A$. Let us forget this result and look at the problem from a different point of view.

If we substitute in the Adler relation (2) the nucleus for the target T then we obtain $\sigma_{\rm tot}(\nu A)\sim\sigma_{\rm tot}(\pi A)$, i.e., the nuclear screening of the neutrino interaction follows directly from the Adler relation. Moreover, there are no restrictions on the energy. The question arises: what is the accuracy of the Adler relation?

Let us write the amplitude $A(\nu T\to lF)$ in the form

$$A(\nu T+lF)=\frac{G}{2^{1/2}}L_\mu M_\mu \tag{4}$$

and separate explicitly in the hadron current M_μ the pion pole contribution:

$$M_\mu=f_\pi\frac{q_\mu}{Q^2+m_\pi^2}A(\pi T\to F)+\tilde{M}_\mu, \tag{5}$$

$\tilde{M}_\mu$ contains the singularities present in the dispersion relation at higher values of Q^2 (the a_1 pole, the $\rho\pi$ cut, etc). It follows from Eq. (5) that demanding conservation of the axial current $q_\mu M_\mu=0$ results in the chiral limit $m_\pi\to 0$ in the relation

$$q_\mu\tilde{M}_\mu=f_\pi A(\pi T\to F). \tag{6}$$

Using this relation in evaluating the cross section at $Q^2=0$ with the help of Eq. (4) leads to the Adler relation.

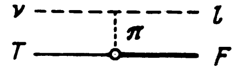

FIG. 3. One-pion exchange diagram for the reaction $\nu T\to lF$.

Above we confined ourselves for convenience to the chiral limit; in what follows this limitation will be removed. Nevertheless we run into a direct contradiction with the preceding result. Indeed, the conclusion that nuclear screening appears only at quite high energies, determined by the a_1 meson mass, is practically independent of the pion mass since in the chiral limit[11] we have $m_a=2\frac{1}{2}m_\rho$, i.e., of the same order as is found experimentally. At the same time it follows from the above derivation that the Adler relation is exact in the chiral limit, i.e., that nuclear screening should exist at any energy.

We thus arrive at a paradox: the neutrino may virtually emit in vacuum only heavy hadronic fluctuations, whose life times are comparable with nuclear dimensions only at quite high energies. On the other hand, it follows from the Adler relation (for $m_\pi=0$) that nuclear screening takes place at any energy. Before resolving this contradiction let us see what is changed if we take $m_\pi\neq 0$. Obviously, if the process is due to distant singularities in the dispersion relation in Q^2 then shifting the pion pole by m_π^2 would lead to small changes in observed quantities on the order of m_π^2/m_a^2. A dispersion relation may also be written in the square of the 4-momentum transferred to the nucleus. In the case of a nucleon target the correction is even smaller, since the quantity $(m_\pi^2+Q^2)/2\nu$ should be equated with the inverse radius of the nucleon. The precision with which the Adler relation is satisfied experimentally[2] in the total cross section for the νN interaction is illustrated in Fig. 4. The contribution of the axial current dominates only at small $Q^2\leqslant 0.1$ (GeV/c)^2, where, unfortunately, the accuracy of the data is poor.

The situation is different in the case of a nuclear target. Here, as is well known,[12] the dispersion relation in the square of 4-momentum transfer has a contribution from anomalous thresholds lying close to zero, at a distance of the order of $1/R_A^2$, where R_A is the root-mean-square radius of the nucleus. Therefore m_π^2 ceases to be a small parameter. The corrections are small only under the condition $(m_\pi^2+Q^2)/2\nu\ll 1/R_A$ (as was already mentioned, in the

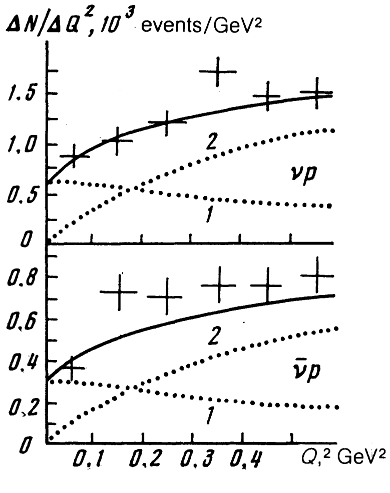

FIG. 4. The Q^2-dependence of the total ν and $\bar{\nu}$ interaction cross section. The experimental points are from Ref. 2. The dotted lines are: 1—calculation with the help of the Adler relation; 2—calculation under the assumption of ρ- and a_1-meson dominance. The solid line is their sum.

case of absorption R_A should be replaced by the mean free path λ). The impression is created that the lifetime of the fluctuation is determined by the pion mass, although we know that the free neutrino cannot virtually emit a pion in vacuum. Thus the contradiction is there for $m_\pi \neq 0$ as well.

The resolution of the paradox consists in the observation that although the neutrino cannot emit a pion in vacuum, it can do so in the nuclear medium as a result of interaction with nucleons. It is precisely the diffractive transitions $\nu N \to l\pi N$ that effectively lead to an increase in the lifetime of the hadronic component of the neutrino in nuclear matter. This is the explanation for the early appearance of nuclear screening at low energies $\nu \geqslant (m_\pi^2 + Q^2)\lambda/2$.

3. INELASTIC CORRECTIONS TO THE TOTAL SCATTERING CROSS SECTION OF NEUTRINOS ON NUCLEI

The above discussion corresponds to the calculation of the cross section for neutrino diffractive scattering on nuclei by the eigenstate method.[13] The neutrino wave function is expanded in the basis of interaction eigenstates:

$$|\nu\rangle = \sum_{\alpha=0}^{\infty} C_\alpha |\nu_\alpha\rangle$$

The highest weight belongs to the interacting component $|\nu_0\rangle$, which describes the "bare" neutrino containing no hadronic fluctuations. For sufficiently high energies mixing may be ignored,[14] i.e., the different components may be viewed as "frozen" for the duration of the interaction. For the calculation of the interaction cross section with the nucleus of an individual component $|\nu_\alpha\rangle$ the eikonal approach is valid. Upon averaging the result with weight $|C_\alpha|^2$ we obtain

$$\sigma_{tot}(\nu A) = 2\int d^2\mathbf{b}\left\{1 - |C_0|^2 - \sum_{\alpha=1}^{\infty} |C_\alpha|^2 \exp[-{}^1\!/_2\sigma_\alpha T(\mathbf{b})]\right\}. \tag{7}$$

Here $\mathbf{b}$ is the impact parameter,

$$T(\mathbf{b}) = \int_{-\infty}^{\infty} dz\rho(\mathbf{b}, z)$$

is a function of the nuclear profile, $\rho(\mathbf{b},z)$ is the nuclear density, and $f_\alpha = i\sigma_\alpha/2$ are the eigenvalues of the scattering amplitude operator. The expression (7) takes into account all inelastic corrections[16] and therefore the method of eigenstates is effectively utilized in the calculation of the cross section for hadron–nucleus interaction. For this, however, one must pay, as a rule, by the appearance of significant model dependence in the quantities σ_α and C_α.

A remarkable peculiarity of the neutrino interactions has to do with the fact that the Adler relation allows one to connect (7) with the cross section for the πA interaction, without knowing σ_α and C_α. If the πA interaction is also expressed in the eigenstate method and compared with (7) then we arrive at the following interesting conclusion: the spectra of the pion eigenstates and of the hadronic component of the neutrino in the axial current, as well as their weights, coincide to within a constant factor.

In the intermediate energy region, where $k_L\pi \geqslant 1/\lambda$ holds, the Adler relation for a nuclear target may be strongly

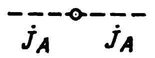

FIG. 5. The diagram corresponding to elastic scattering of the axial current off a nucleon in the nucleus.

violated. The eigenstates method also becomes ineffective due to strong mixing, which is difficult to take properly into account. In this energy region the approach of Glauber[15] and Grivob[16] works best. We consider the contribution to the total cross section of the diagrams in Figs. 5 and 6 (in accordance with the optical theorem we calculate instead of the total cross section the elastic forward scattering amplitude). The first diagram corresponds to the interaction of the neutrino with one of the nucleons in the nucleus. This contribution has the form

$$\sigma_{tot}^{(1)}(\nu A) = A\sigma_{tot}(\nu N). \tag{8}$$

The contribution of the second diagram (Fig. 6) is the inelastic correction to the first, corresponding to the pion production in the intermediate state. It is calculated from a formula analogous to the Karmanov–Kondratyuk formula:[17]

$$\sigma_{tot}^{(2)} = 2f_{\nu\pi}^2 \int d^2\mathbf{b} \int_{-\infty}^{\infty} dz_1 \rho(\mathbf{b}, z_1) \int_{-\infty}^{z_1} dz_2 \rho(\mathbf{b}, z_2)$$
$$\times \cos[k_L^\pi(z_1 - z_2)]\exp\left[-f_{\pi\pi}\int_{z_1}^{z_2} dz\rho(\mathbf{b}, z)\right]. \tag{9}$$

Here $f_{\pi\pi}$ and $f_{\nu\pi}$ are the imaginary parts of the amplitudes for elastic πN scattering and for the process of diffractive neutrino-production of pions on nucleons, normalized so that $f_{\pi\pi} = \sigma_{\rm tot}/2$ holds; z is the longitudinal coordinate, $\rho(\mathbf{b},z)$ is the nuclear density of nucleons, and $k_L^\pi = (m_\pi^2 + Q^2)/2\nu$ is the longitudinal momentum transfer in the process of neutrino-production of the pion, which determines the phase shift $k_L\pi(z_1 - z_2)$. The exponential factor in Eq. (9) takes into account the requirement that the pion passes through the nucleus without interacting.

It is interesting to trace the cancellation of the volume terms in Eqs. (8) and (9) at asymptotic energies, when $k_L\pi \ll 1$ holds. Under that condition Eq. (9) may be rewritten as

$$\sigma_{tot}^{(2)}|_{k_L\to 0} = -\left(\frac{f_{\nu\pi}}{f_{\pi\pi}}\right)^2 [A\sigma_{tot}(\pi N) - \sigma_{tot}^0(\pi A)], \tag{10}$$

where

$$\sigma_{tot}^0(\pi A) = 2\int d^2\mathbf{b}\{1 - \exp[-f_{\pi\pi}T(\mathbf{b})]\} \tag{11}$$

is the total cross section for the pion–nucleus interaction, calculated in the eikonal approximation.

It is seen from a comparison of Eqs. (8) and (10) that in their sum the volume terms, proportional to A, cancel provided that

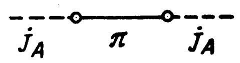

FIG. 6. The diagram describing pion production in the intermediate state.

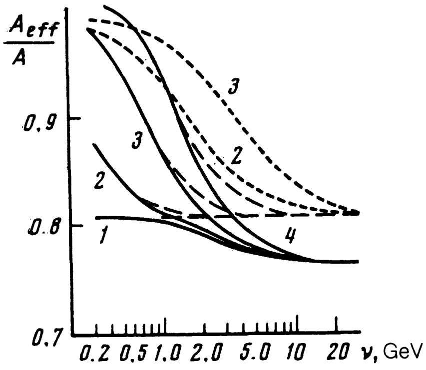

FIG. 7. The A_{eff}/A ratio for neutrino interactions with Ne nuclei as a function of the neutrino energy at various values of Q^2. (Curve 1 is for $Q^2 = 0$ GeV^2, curve 2 for 0.2 GeV^2, curve 3 for 0.5 GeV^2, and curve 4 for 1.0 GeV^2). The solid and dashed lines show the results of calculations respectively with and without diffraction excitation of the pion in the intermediate state (the Glauber–Gribov model). The dotted lines correspond to Bell's optical model.

$$\frac{\sigma_{tot}(\nu N)}{\sigma_{tot}(\pi N)} = \left(\frac{f_{\nu\pi}}{f_{\pi\pi}}\right)^2. \qquad (12)$$

But that is a direct consequence of the Adler relation (2), where it is necessary to sum over all final states F.

Thus a violation of the Adler relation would manifest itself in the present case in the appearance of the volume term. The experimental study of the A-dependence of the cross section $\sigma_{tot}(\nu A)$ provides yet another interesting test of the Adler relation, and therefore a test of PCAC.

We note that the condition (12) does not depend on the size of the factor standing in front of $\sigma(\pi T \to F)$ on the right side of Eq. (2). It is only necessary that it be independent of the final state F. In contrast to this the comparison with the data shown in Fig. 4, and the processes of coherent neutrino-production of hadrons on nuclei test the size of the coefficient in Eq. (2).

In the intermediate energy region the quantity k_L^π cannot be neglected and the volume terms in Eqs. (8) and (9) do not cancel. We were not able to evaluate (9) analytically and the results of numerical calculations of the ratio $1/A\sigma_{tot}(\nu A)/\sigma_{tot}(\nu N)$ for the Ne nucleus are shown in Fig. 7. For the nuclear density $\rho(\mathbf{r})$ we used the Woods–Saxon parametrization:

$$\rho(\mathbf{r}) = \rho_0\left[1 + \exp\left(\frac{r-R}{a}\right)\right]^{-1},$$

$$\rho_0 = \frac{3A}{4\pi R^3}\left(1 + \frac{\pi^2 a^2}{R^2}\right)^{-1}$$

with the parameters R and a from Ref. 18. It is seen from Fig. 7 that at $Q^2 = 0$ the nuclear screening is practically saturated even at energies of a few hundred MeV. With increasing Q^2 the passage to the asymptotic regime, shown in Fig. 7 by broken traces, naturally tightens.

It is of interest to compare our calculations with the results obtained in Bell's optical model,[3,9] since this model is used[1,2] in the interpretation of the experimental data. The optical model equation, which describes the dependence on the longitudinal coordinate z of the probability amplitude to detect the hadronic fluctuation of the neutrino in nuclear matter, has the form[3,9]

$$\frac{d\Psi}{dz} = -\frac{1}{\lambda}\Psi + \frac{1}{\lambda}\frac{q_L^\pi}{q_L^\pi - if_{\pi\pi}\rho(z)}\Psi_0. \qquad (13)$$

After solving this equation and averaging Ψ over the volume of the nucleus we obtain the following expression for $\sigma_{tot}(\nu A)$:

$$\sigma_{tot}^{(OM)}(\nu A) = \left(\frac{f_{\nu\pi}}{f_{\pi\pi}}\right)^2 \sigma_{tot}^0(\pi A) + 2(f_{\nu\pi})^2 \int d^2\mathbf{b} \int_{-\infty}^{\infty} dz_1 \rho(\mathbf{b}, z_1)$$
$$\times \int_{-\infty}^{z_1} dz_2 \rho(\mathbf{b}, z_2) \frac{q_L^\pi}{q_L^\pi - if_{\pi\pi}\rho(z)} \exp\left[-f_{\pi\pi}\int_{z_1}^{z_2} dz\rho(\mathbf{b}, z)\right]. \qquad (14)$$

This expression, evidently, does not coincide with Eqs. (8) and (9). A comparison between our results and the results of numerical computations, carried out at $Q^2 = 0.2$ and 0.5 (GeV/c)^2, is shown in Fig. 7. It is seen that the difference is large. The optical model noticeably underestimates the size of nuclear screening in the intermediate region of energies. The asymptotic behavior is the same in both approaches, a result which is trivial as it is dictated by the Adler relation.

We consider now more complicated inelastic corrections connected with the possibility of diffraction dissociation of the pion. The corresponding diagrams are shown in Fig. 8. The first (Fig. 8a) corresponds to the summary contribution of processes of diffractive neutrino-production of all possible states other than the pion: $a_1, \rho\pi, 3\pi, \ldots$, denoted in the figure by X. The diagrams in Fig. 8b and 8c, take into account the possibility of diffractive transitions of the pion into these states.

It is not hard to see that the contribution of each of the three diagrams in Fig. 8 contains the volume term. With the Adler relation taken into account the volume terms cancel at asymptotic energies. Moreover, independently of the size of ν the contributions of the diagrams in Fig. 8a and 8b are completely cancelled by analogous terms contained in the expression for the contribution of the diagram in Fig. 8c. Leaving out some simple steps we write the expression for the summary contribution of the diagrams in Fig. 8 as:

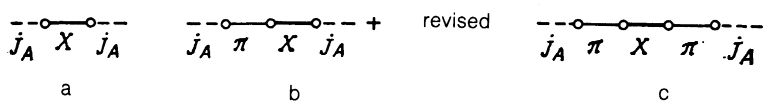

FIG. 8. Diagrams that take into account processes of diffraction excitation of the axial current and the pion in interaction with nucleons in the nucleus.

$$\sigma_{tot}^{(3)}(\nu A) = -4\pi\left(\frac{f_{\nu\pi}}{f_{\pi\pi}}\right)^2 \int d^2\mathbf{b} \exp[-f_{\pi\pi}T(\mathbf{b})] \int_{-\infty}^{\infty} dz_1 \rho(\mathbf{b}, z_1)$$

$$\times \int_{-\infty}^{z_1} dz_2 \rho(\mathbf{b}, z_2) \int_{9m_\pi^2}^{\infty} dM_x^2 \frac{d\sigma_{DD}(\pi N)}{dM_x^2\, dk_T^2}\bigg|_{k^2_T=0} \cos[k_L^x(z_1 - z_2)]. \tag{15}$$

We make here the simplifying assumption that for all the states that are diffractive excitations of the pion one has $\sigma_{\rm tot}(XN) = \sigma_{\rm tot}(\pi N)$. Further, since we have $k_L^x = (m_x^2 + Q^2)/2\nu \gg k_L^\pi$, we assume that k_L^π may be set equal to zero because the correction (15) becomes noticeable at quite high energies only.

Equation (15), accurate to within a factor $(f_{\nu\pi}/f_{\pi\pi})$,[2] coincides with the Karmanov–Kondratyuk formula,[17] and therefore the addition of $\sigma_{\rm tot}^{(3)}(\nu A)$ is equivalent to taking into account in the surface term $\sigma_{\rm tot}^{(2)}(\nu A)$ of the inelastic correction to $\sigma_{\rm tot}(\pi A)$. At asymptotic energies this is trivial, being a consequence of the Adler relation. Nevertheless it turns out to be true also at intermediate energies. By using a general technique, developed by Bertocchi and Treleani[19] for the vector current, we can show that this conclusion is preserved also after taking into account higher order inelastic corrections.

To arrive at a numerical estimate of the correction (15) we take into account the fact that the mass spectrum in the process of diffraction dissociation is mostly centered in the region of the a_1-meson mass. Therefore in the expression for k_L^x we set $M_x = m_a$. The total cross section for the forward diffraction dissociation of the pion may be estimated by taking 2/3 of the cross section for the process $pp \to pX$, which yields about 10 mbn/GeV.[2] After taking into account the contribution of $\sigma_{\rm tot}^{(3)}(\nu A)$ the quantity $A_{\rm eff}$ for the Ne nucleus is lowered in the asymptotic region by approximately 10%. The corresponding ν-dependence for various values of Q^2 is shown in Fig. 7 by the solid lines.

Finally, a comparison can be carried out with experimental data.[1,2] To this end it is necessary to average $\sigma_{\rm tot}(\nu A)$ over the internal $0 \leqslant Q^2 \leqslant 0.2$ (GeV/c)^2. The dependence on Q^2 is contained, besides the momentum transfer k_L, in the form factor

$$F(Q^2) = m_a^2/(Q^2 + m_a^2).$$

This dependence does not necessarily presume a_1-dominance of the axial current. As was shown in Ref. 20, the Q^2-dependence corresponding to the $\rho\pi$ cut imitates the a_1 pole to high accuracy.

It is also necessary to take into account the contribution of the vector current, which on the assumption of ρ-dominance is given by the sum of Eqs. (8) and (9), where one should replace $f_{\nu\pi}$ by $f_{\nu\rho}$, $f_{\pi\pi}$ by $f_{\rho\rho}$, k_L^π by $k_L\rho$, i.e., m_π by m_ρ. Further, we have[10]

$$\left(\frac{f_{\nu\rho}}{f_{\rho\rho}}\right)^2 = \frac{G^2}{2\pi^2}\frac{1}{\nu} f_\rho^2 Q^2 (1 - y + y^2/2)(m_\rho^2 + Q^2)^{-2}, \tag{16}$$

where $y = \nu/E$, $f_\rho = 2\frac{1}{2}m_\rho^2/\gamma_\rho$, where γ_ρ is the universal hadron coupling constant ($\gamma_\rho^2/4\pi \approx 2.4$) Due to the conservation of the vector current its contribution, as can be seen from Eq. (16), vanishes for $Q^2 = 0$. The relation between the contributions of the vector and axial currents to the neutrino–nucleon interaction cross section was already shown in Fig. 4. In the case of a nuclear target these contributions are calculated with the help of the procedure described above. As a result of averaging the cross section over Q^2 in the interval $0 \leqslant Q^2 \leqslant 0.2$ (GeV/c)^2 we obtain a ν-dependent screening factor $A_{\rm eff}/A$, as is shown for the Ne nucleus in Fig. 9. The calculation agrees well with experimental data.[1,2]

Such a comparison provides, in principle, a test of PCAC, since the coefficient of the volume term vanishes only if the Adler relation is satisfied (for sufficiently high energies). Although the PCAC-violating contribution is relatively strengthened by a factor $A^{2/3}$, the present accuracy of the data of Ref. 1 and 2 allow for significant violation of the Adler relation.

4. CONCLUSION

The partial conservation of the axial current phenomenon is a reflection of the presence in the theory of an initial symmetry of the strong interactions with respect to chiral transformations and the presence in nature of a small parameter—the square of the pion mass. While the transversality of the axial currents of massless quarks is a trivial property, the reasons for the conservation of the axial current in the sector of hadrons, which have acquired large masses as a result of spontaneous breaking of chiral symmetry, are as yet not understood. The PCAC conditions allows us in a remarkable fashion to relate quantities which, at first glance, have nothing in common. For example, it relates the pion-nucleon coupling constant and the $\pi \to \mu\nu$ decay constant (the Goldberger–Treiman relation). It should be emphasized that each new manifestation of PCAC permits an independent experimental test of this phenomenon, since the structure of the hadronic currents differs in different processes and the conservation of the current in one case provides no guarantees against a strong violation in another case.

Up till now the most careful test of PCAC consequences were carried out in low-energy pion physics. Recently new possibilities have arisen in high-energy physics. For example, the lepton decays $\tau \to 3\pi + \nu_\tau$ give direct information on

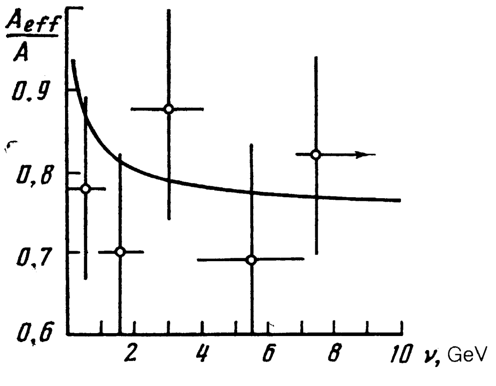

FIG. 9. The $A_{\rm eff}/A$ ratio for the Ne nucleus for $x < 0.2$ and $Q^2 < 0.2$ (GeV/c)^2. The curve shows results of a calculation that includes the contributions of the axial and vector currents.

the spectral function of the axial current. The recently obtained experimental data[12] allowed for the first time for a direct test of the Weinberg sum rules, and have verified them to within experimental errors. The measured distribution in the effective 3π mass as yet does not contradict the a_1-dominance hypothesis and the smallness of the longitudinal component of the spectral function (which follows from PCAC). It is necessary, however, to carry out a partial-wave analysis of the 3π system to verify these assertions.

In the present work we have considered a special part of high-energy physics—the diffractive scattering of the axial current. The study of such processes has become possible with the appearance of intense high-energy neutrino beams. The unique data obtained by the WA59 collaboration have proven the existence of nuclear screening for the neutrino interaction. Although these data, shown in Fig. 9, agree with theoretical calculations it is too soon to speak of a new experimental confirmation of PCAC. To this end, as was demonstrated, data of much higher precision are needed.

Favorable conditions for the extraction of the contribution of the axial current are provided from this point of view by the study of coherent neutrino-production of hadrons on nuclei. In these reactions nondiffractive mechanisms, connected with a change in G-parity, are forbidden and, consequently, the contribution of the vector current is suppressed. This explains the heightened interest in experiments on coherent processes. Unfortunately (or fortunately?) the agreement of the calculations in Ref. 20, based on consequences of PCAC, with the data on coherent neutrino-production of π and 3π on Ne nuclei can hardly be called striking. Nevertheless it is too soon to speak of a serious disagreement. It is necessary to both improve the precision of the experimental data and to carry out model-independent calculations.

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[1] BEBC WA59 Collaboration, P. P. Allport *et al.*, Paper submitted to the XIX Int. Symp. on Multiparticle Dynamics, Arles, France, 1988.
[2] BEBC WA59 Collaboration, P. Marage, Report submitted to the XIII Int. Conf. on Neutrino Physics and Astrophysics, Boston, 1988.
[3] J. S. Bell, Phys. Rev. Lett. **13**, 57 (1964).
[4] M. Ross and L. Stodolsky, Phys. Rev. **149**, 1172 (1966).
[5] S. D. Drell and S. Trefil, Phys. Rev. Lett. **16**, 552, 832 (1966).
[6] T. H. Bauer, R. D. Spital, D. R. Yennie, and F. M. Pipkin, Rev. Mod. Phys. **50**, 261 (1978).
[7] B. Z. Kopeliovich, Phys. Lett. **227B**, 461 (1989).
[8] S. Adler, Phys. Rev. **135**, 963 (1964).
[9] J. Bell, *Lectures in Particles and Fields*, edited by H. H. Aly, Gordon and Breach, New York, 1969, p. 323.
[10] C. Piketty and L. Stodolsky, Nucl. Phys. **B15**, 571 (1970).
[11] S. Weinberg, Phys. Rev. Lett. **18**, 507 (1968).
[12] R. Karplus, C. M. Sommerfield, and E. H. Wichmann, Phys. Rev. **111**, 1187 (1958).
[13] B. Z. Kopeliovich and L. I. Lapidus, Pis'ma Zh. Eksp. Teor. Fiz. **28**, 664 (1978) [JETP Lett. **28**, 614 (1978)]; *Proc. of the V International Seminar on Multiple Production of Hadrons at High Energies*, Dubna, 1978, p. 460.
[14] B. Z. Kopeliovich and L. I. Lapidus, Pis'ma Zh. Eksp. Teor. Fiz. **32**, 612 (1980) [JETP Lett. **32**, 597 (1980)].
[15] R. J. Glauber, *Lectures in Theoretical Physics*, 1958, Vol. 1, edited by W. E. Brittin and L. G. Dunham, New York, 1959.
[16] V. N. Gribov, Zh. Eksp. Teor. Fiz. **56**, 892 (1969) [Sov. Phys. JETP **29**, 483 (1969)].
[17] V. Karmanov and L. A. Kondratyuk, Pis'ma Zh. Eksp. Teor. Fiz. **18**, 266 (1973) [JETP Lett. **18**, 451 (1973)].
[18] H. De Vries, C. W. De Jager, and C. De Vries, Atomic Data and Nuclear Data Tables **36**, 495 (1987).
[19] L. Bertocchi and D. Treleani, Nuovo Cimento **A54**, 113 (1979).
[20] A. A. Bel'kov and B. Z. Kopeliovich, Yad. Fiz. **46**, 874 (1987) [Sov. J. Nucl. Phys. **46**, 499 (1987)].
[21] ARGUS Collaboration, H. Albrecht *et al.*, Z. Phys. C **33**, 7 (1986).

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