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ISOVECTOR KAON FORM FACTOR IN THE DISPERSION THEORY AND S-WAVE πK-SCATTERING LENGTHS: Preprint ITEP 93-29/

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A one-dimensional integral representation for isovector kaon form factor is constructed within the dispersion theory in terms of the pion form factor and the backward πK-scattering amplitude. The normalization condition for isovector kaon form factor at zero momentum transfer gives a sum rule for the πK-scattering amplitude, with the use of which difference between the s-wave πK-scattering lengths in triplet and singlet isospin states is estimated to be $a_0^{3/2} - a_0^{1/2} \approx 0.22\,\mu^{-1}$ where μ is the pion mass. In agreement with the vector-meson-dominance model, deviations of the isovector kaon form factor from half of the pion form factor are found to be small.

Fig. - 2, ref. -

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I. INTRODUCTION

The experimental studying the pion, kaon, and nucleon form factors [1-4] stimulates development of phenemenological models capable of the detailed description of the hadron form factors at low and intermediate energies [5-7] and having correct QCD asymptotic behavior. The dispersion theory of strong interactions based on the unitarity, crossing symmetry, and analyticity is known to be useful in analyzing the scattering amplitudes and the form factors. In this way, the ρ-meson resonance was first predicted [8] by studying the nucleon form factor data.

In this paper we consider the problem of isovector kaon form factor in the dispersion theory. A part this work, connected with properties of the pion form factor, rests essentially on the Frazer-Fulco-Gounaris-Sakurai (FFGS) model [8,9], that has correct analytical properties, satisfies the low-energy unitarity and crossing symmetry. This model gives a simple relation between imaginary part of the pion form factor and absolute square of the pion form factor, as a result of which the imaginary part of the kaon form factor determined by the unitarity condition turns out to be proportional to the imaginary part of the pion form factor. The dispersion relation for the kaon form factor can then be used to express the kaon form factor in terms of the pion form factor and the backward πK-scattering amplitude in one-dimensional integral form. Using this representation, we give simple estimates of the kaon form factor and discuss relation to the vector-meson-dominance model. The asymptotics of the kaon form factor at high momenta transfer is determined by asymptotics of the pion form factor, which has within the FFGS model a $O(1/t)$ behavior at momenta $|t| < 10^6$ GeV^2. Our results make sense in this region. The asymptotics of the kaon form factor is in agreement with the quark counting rules.

The paper is organized as follows. In Sect. II we review basic properties of the πK-scattering amplitude, including a backward dispersion relation used in computations of the kaon form factor, and discuss briefly the FFGS model for the pion form factor. In Sect. III, the unitarity relation for the kaon form factor is derived and a simple estimate of the kaon form factor, based on the assumption that ratio between the πK-scattering amplitude and the pion form factor

represents a slowly variable function, is given. This assumption is equivalent in fact to the ρ-meson dominance in the πK-scattering amplitude. In this approximation, the isovector kaon form factor coincides with the pion form factor. This result is exact in the $SU(3)_f$ limit and also in the no-width limit of the ρ-meson. In Sect. IV a one-dimensional integral representation for the kaon form factor in terms of the pion form factor and the backward πK-scattering amplitude of the s-channel is derived and a sum rule for the amplitude following from the normalization condition for the kaon form factor is constructed. In Sect. V the sum rule is saturated by lowest partial amplitudes to estimate difference between the s-wave πK- scattering lengths in triplet and singlet isotopic channels. The result is in agreement with the experimental data and predictions of the effective chiral theory. The kaon form factor is evaluated in the same way and found to be close numerically to half of the pion form factor. In Sect. VI we discuss the results.

II. THE πK-SCATTERING AMPLITUDE AND THE FFGS MODEL FOR PION FORM FACTOR

The unitarity relation expresses the imaginary part of the kaon isovector form factor through the πK- scattering amplitude and the pion form factor. The analytical properties of the amplitude and a dispersion relation which we use in this work are discussed in ref.[10] on the basis of the Mandelstam double representation. The problem of isovector kaon form factor has many features in common and can be useful in studying isovector nucleon form factors [11,12]. The distinction is due to spin of the nucleons and absence of the pole Born terms.

The πK- scattering amplitude is defined by

$$i(2\pi)^4\delta(q_2+k_2+q_1+k_1)\mathfrak{I}(q_2,k_2,q_1,k_1)_{j,\beta,i,\alpha}=$$

$$\frac{1}{Z_\pi Z_K}\int dx_2dx_1dy_2dy_1 e^{-iq_2x_2-ik_2y_2-iq_1x_1-ik_1y_1}(\Box_{x_2}+\mu^2)(\Box_{y_2}+m^2)(\Box_{x_1}+\mu^2)(\Box_{y_1}+m^2) \qquad (2.1)$$

$$<0|T(\pi_j(x_2)\pi_i(x_1)\phi_\beta(y_2)\phi^+_\alpha(y_1))|0>.$$

where Z_π and Z_K are the renormalization constants of the pion and kaon propagators, $\pi_i(x)$ and $\phi_\alpha(x)$ the pion and kaon fields. The components of the field $\phi_\alpha(x)$ are identified with K-mesons by

$$\phi_\alpha = \begin{Bmatrix} K^+ \\ K^0 \end{Bmatrix}. \tag{2.2}$$

The isospin operator transforms the pion wave function as $\omega_i\pi_j = -i\varepsilon_{ijk}\pi_k$.

There are two invariant amplitudes in terms of which the isospin content of the processes can be expressed. These amplitudes, $A(s,u,t)$ and $B(s,u,t)$, are defined by

$$\mathfrak{I}(p_2,k_2,p_1,k_1)_{i,\beta,j,\alpha} = \delta_{ji}\delta_{\beta\alpha}A(s,u,t) + i\varepsilon_{kji}\tau^k_{\beta\alpha}B(s,u,t), \tag{2.3}$$

where $s = (p_1 + k_1)^2$, $t = (p_1 + p_2)^2$, $u = (p_1 + k_2)^2$, and τ are Pauli matrices.

Crossing symmetry and physical equivalence of the channels connected by permutations $q_1 \leftrightarrow -q_2$, $i \leftrightarrow j$ implies that $A(s,u,t) = A(u,s,t)$ and $B(s,u,t) = -B(u,s,t)$. The amplitudes A and B correspond to superpositions of states with definite total isospins. The decomposition in the s-, u-, and t- channels can be done in the form

$$\mathfrak{I} = \sum_{I=1/2,3/2} A^s_I P_I = \sum_{I=1/2,3/2} A^u_I P_I = \sum_{I=0,1} A^t_I P_I \;, \tag{2.4}$$

where P_I are projection operators with definite isotopic spin. The isotopic amplitudes can be found to be $A^s_{1/2} = A - 2B$, $A^s_{3/2} = A + B$, $A^u_{1/2} = A + 2B$, $A^u_{3/2} = A - B$, $A^t_0 = A/\sqrt{6}$, $A^t_1 = 2B$.

We use in the next section a backward dispersion relation for the function

$$\beta(s,u,t) = \frac{B(s,u,t)}{s-u} \tag{2.5}$$

which is regular at $s = u$ and symmetric under the substitution $s \leftrightarrow u$.

The momentum squared in the center-of-mass frame of the s-channel, $\nu = \lambda(s,m^2,\mu^2)/(4s)$, where $\lambda(s,m^2,\mu^2) = (s-(m+\mu)^2)(s-(m-\mu^2))$, is treated as argument of the amplitude. The s- channel scattering angle is defined by $\cos\theta_s = 1 - 2st/\lambda(s,m^2,\mu^2)$. Analytical properties of the function $\beta(s,u,t)$ as a function of the variable ν at fixed scattering angle θ_s are determined by

analytical properties of the amplitude $B(s,u,t)$ with respect to the variable s, which follow in turn from the Mandelstam double representation. In case of the backward scattering [10], the function $\beta(\nu,\cos\theta_s = -1)$ is analytical in the complex ν-plane with two cuts $(-\infty,-\mu^2)$ and $(0,+\infty)$.

We shall use the Frazer-Fulco-Gounaris-Sakurai model of the pion form factor [8,9] where

$$F_\pi(t) = \frac{D(0)}{D(t)}, \tag{2.6}$$

$D(t) = a + bt + k^2(t)h(t)$, a and b are coefficients fixed by fitting the ρ-meson mass and width, and $h(t)$ is given by

$$h(t) = \frac{1}{16\pi^2}\sqrt{\frac{4\mu^2 - t}{-t}}\ln\left(\frac{\sqrt{4\mu^2 - t} + \sqrt{-t}}{\sqrt{4\mu^2 - t} - \sqrt{-t}}\right). \tag{2.7}$$

Notice that $\mathrm{Im}D(t+i0) = -k(t)^3/(8\pi\sqrt{t})$, so that

$$\frac{k(t)^3}{8\pi\sqrt{t}}|F_\pi(t)|^2 = \mathrm{Im}F_\pi(t+i0)D(0). \tag{2.8}$$

III. UNITARITY, ELIMINATION OF THE LEFT CUT, AND RELATION TO THE VMD MODEL

The isoscalar and isovector kaon form factors $F_{s,v}(t)$ are defined by

$$\begin{aligned}(2\pi)^4\delta(q+k_2+k_1)(k_1-k_2)_\mu(F_s(t+i0)\delta_{\beta\alpha}+F_v(t+i0)\tau^3_{\beta\alpha}) = \\ -\frac{1}{Z_K}\int dx\,dy_2\,dy_1\,e^{-iqx-ik_2y_2-ik_1y_1}(\partial_{y_2}+m^2)(\partial_{y_1}+m^2)\langle 0|T(j_\mu(x)\phi_\beta(y_2)\phi^+_\alpha(y_1))|0\rangle.\end{aligned} \tag{3.1}$$

The kaon charge operator has the form $e_{\alpha\beta} = (\delta_{\alpha\beta} + \tau^3_{\alpha\beta})/2$. The form factors are normalized by $F_{s,v}(0) = 1/2$.

The unitarity relation for two-kaon state $\langle f| = \langle k_1,\alpha;\, k_2,\beta;\, \mathrm{out}|$ gives

$$\langle f\ \mathrm{out}|j_\mu(q)|0\rangle - \langle f\ \mathrm{in}|j_\mu(q)|0\rangle = i\sum_n(2\pi)^4\delta(P_f+P_n)\mathcal{T}_{fn}\langle n\ \mathrm{in}|j_\mu(q)|0\rangle. \tag{3.2}$$

Taking into account two-pion intermediate states we get

$$\mathrm{Im}F_K^v(t) = \frac{t-4\mu^2}{4\sqrt{t-4m^2}}\, h_1^1(t) F_\pi(t)^*. \tag{3.3}$$

The $l = 1$ t- channel partial wave amplitudes are defined according to the decomposition

$$A_l^1 = 2B = 8\pi\sqrt{t}\sum_l (2l+1) h_l^1 P_l(\cos\theta_t) \tag{3.4}$$

In case of the backward πK- scattering $\cos\theta_s = \cos\theta_t = -1$, $t = -4v$, $4K(v) = -((t-4m^2)(t-4\mu^2))^{1/2}$, and $\beta(s,u,t) = B(s,u,t)/(4K(v))$. In the vicinity of $v = -\mu^2$ we neglect by higher partial waves ($l = 3,5,\ldots$) and get

$$h_1^1(v) = \frac{K(v)}{3\pi\sqrt{t}}\beta(v,-1) \tag{3.5}$$

and

$$\mathrm{Im}F_K^v(t) = \frac{(t-4\mu^2)^{3/2}}{48\pi t}\beta(v,-1)F_\pi(t)^*. \tag{3.6}$$

The ratio $\beta(v,-1)/F_\pi(v)$ is regular on the left cut $(-\infty, -\mu^2)$, since we neglect by the higher partial waves. The pion form factor in the FFGS model has no zeros on the complex v-plane, so we can write for $\beta(v,-1)/F_\pi(v)$ a one-subtracted dispersion relation

$$\frac{\beta(v,-1)}{F_\pi(v)} = \frac{\beta(0,-1)}{F_\pi(0)} + \frac{v}{\pi}\int_0^{+\infty}\frac{\mathrm{Im}\beta(v',-1)}{v'(v'-v)}\frac{dv'}{F_\pi(v')}. \tag{3.7}$$

The imaginary part of the kaon form factor with the use of eqs.(2.8) and (3.6) becomes

$$\mathrm{Im}F_K^v(v) = \frac{4D(0)}{3}\mathrm{Im}F_\pi(v)\left(\frac{\beta(0,-1)}{F_\pi(0)} + \frac{v}{\pi}\int_0^{+\infty}\frac{\mathrm{Im}\beta(v',-1)}{v'(v'-v)}\frac{dv'}{F_\pi(v')}\right). \tag{3.8}$$

The amplitude at zero momentum is expressible in terms of the s-wave scattering lengths

$$\beta(0,-1) = \frac{2\pi(m+\mu)}{3m\mu}\left(-a_0^{1/2} + a_0^{3/2}\right) \tag{3.9}$$

In physical region of the s-channel, the value $\mathrm{Im}\beta(0,-1)$ has the form

$$\mathrm{Im}\beta(\nu,-1) = \frac{2\pi\sqrt{s}}{3\sqrt{\nu}\,K(\nu)} \sum_l (2l+1)(-)^l (\sin^2(\delta_l^{1/2}(\nu)) - \sin^2(\delta_l^{3/2}(\nu))). \qquad (3.10)$$

A suggestion that the ratio $\beta(\nu,-1)/F_\pi(\nu)$ represents a slowly variable function is equivalent to a hypothesis of the ρ-meson dominance in the πK-scattering amplitude, since the pion form factor is determined by the ρ-meson pole. In this approximation we get

$$\mathrm{Im}F_K^\nu(\nu) = C\,\mathrm{Im}F_\pi(\nu) \qquad (3.11)$$

The general solution has the form

$$F_K^\nu(t) = CF_\pi(t) + C_1 + C_2 t + C_3 t^2 + \ldots \qquad (3.12)$$

The asymptotic behavior of the pion and kaon form factors are identical, so $C_1 = C_2 = \ldots = 0$. Using normalization condition $C = 1/2$, we get

$$F_K^\nu(t) = \frac{1}{2} F_\pi(t). \qquad (3.13)$$

In the vector-meson-dominance model, this result is exact (up to uncertainties connected with unknown value of the ρKK form factor). It satisfies also the $SU(3)_f$ symmetry, despite the unitarity conditions for the pion and kaon form factors are different.

IV. ONE-DIMENSIONAL INTEGRAL REPRESENTATION AND SUM RULE

The imaginary part for the kaon form factor given by eq.(3.8) can be used to reconstruct the kaon form factor through the dispersion relation. The dispersion integral can be evaluated explicitly to give

$$F_K^\nu(\nu) = \frac{4D(0)}{3} F_\pi(\nu)\left(\frac{\beta(0,-1)}{F_\pi(0)} - \frac{1}{\pi}\int_0^{+\infty} d\nu' \frac{\mathrm{Im}\beta(\nu',-1)}{\nu'} \frac{\nu/F_\pi(\nu') - \nu'/F_\pi(\nu)}{(\nu-\nu')}\right). \qquad (4.1)$$

The integral extends over the crossing s-channel. It is possible to add a polinom to the right side of this equation, as in eq.(3.13). The kaon form factor, however, is proportional to the pion form factor and has correct asymptotic behavior, so such a polinom should be set equal to zero.

Substituting $F_\pi(t) = (1 - t/m_\rho^2)^{-1}$ in eq.(4.1) we find that the integral term does not depend on the v. Therefore eq.(3.13) is exact in the no-width limit of the ρ-meson and corrections to eq.(3.13) are of order $O(\Gamma/m)$.

The kaon form factor in the dispersion theory has normally a two-dimensional spectral representation, the first integral comes from the dispersion relation for the kaon form factor, while the second one comes from the dispersion relation for the amplitude through eq.(3.7). Using eq.(3.3) and analytical properties of the integrand in the complex v-plane, it becomes possible to evaluate one integral explicitly, as a result of which we get the one-dimensional integral representation (4.1).

The normalization condition at zero momentum transfer gives a sum rule for the backward scattering amplitude

$$\frac{1}{2} = \frac{4D(t)}{3}\left(\beta(t,-1) - \frac{1}{\pi}\int_0^{1-m} dv' \frac{\mathrm{Im}\beta(v',-1)}{v'}\right). \tag{4.2}$$

The dependence on the pion form factor enters through the value D(t) only.

The approximation discussed in previous Sect. is equivalent to assumption that the first term in eq.(4.2) determined by the s-wave πK-scattering lengths is dominant. We wish to exploit now an opposite possibility, namely that the integral term is dominant as compared to the scattering lengths. To give an idea about possible effect, we evaluate imaginary part of the scattering amplitude in the delta function approximation: $\mathrm{Im}\beta(v,-1) = \gamma\delta(v - v_0)$. The constant γ can be absorbed by redefinition of the normalization constant. The result becomes

$$F^v_K(t) = F_\pi(t)\frac{tF_\pi(t_0) - t_0F_\pi(t)}{2(t - t_0)}. \tag{4.3}$$

In the no-width limit we reproduce eq.(3.13). If we take into account in the same approximation only phase shift generated by strange K_V^*-meson, by writing $\sin^2(\delta_1^{1/2}(\nu)) = \pi\Gamma_*K(\nu_*)/m_*\delta(\nu - \nu_*)$, we obtain $t_* = -4m_* = -0.33\ GeV^2$, $F_\pi(t_*) = 0.65$. The absolute value of the form factor (4.3) for $t_* = -0.33\ GeV^2$ is shown in Fig.1 (curve B).

V. NUMERICAL RESULTS

We saturate the imaginary part of the backward scattering amplitude by lowest partial waves $l = 0$ with isospin $I = 1/2, 3/2$ and $l = 1$ with isospin $I = 1/2$. The $l = 1$ partial wave with isospin $I = 3/2$ is small and can be neglected. The experimental phase shifts [14] in the $l = 0$ wave are interpolated smoothly by curves with $O(k)$ limiting behavior as $k \to 0$ and $\pi + O(1/k)$ behavior as $k \to \infty$. The $l = 1$ partial wave with isospin $I = 1/2$ is interpolated by the Breit-Wigner amplitude with the K_V^*-meson momentum-dependent width with limiting behavior $O(k^3)$ as $k \to 0$ and $O(k)$ as $k \to \infty$, according to the effective range expression, so that the amplitude has an $O(1/k)$ order of magnitude as $k \to \infty$.

The numerical results can be summarized as follows

$$(0.50) = 2.51\ \mu(a_0^{3/2} - a_0^{1/2}) + (-0.33 + 0.06 + 0.22 + \ldots) \qquad (5.1)$$

In the second term we display contributions of the partial waves $(l,I) = (0,1/2), (0,3/2), (1,1/2)$. The contribution of the K_V^*-meson can be estimated with good precision in the delta function approximation. The unitarity condition with two pions in the intermediate state makes sense at low energies only, so taking into account higher partial waves is beyond accuracy of our approach. We can check, however, consistency of the sum rule by estimating contributions from more massive resonances coupled to the πK-channel. The resonances in partial waves $(l,I) = (0,1/2), (1,1/2), (1,1/2), (2,1/2), (3,1/2), (4,1/2)$ with masses, respectively, 1.43, 1.37, 1.68, 1.43, 1.77, 2.05 GeV contribute to the sum rule as $-0.05 + 0.01 + 0.04 - 0.04 + 0.02 - 0.01 = -0.03$. The resonance $(1,1/2)$ with mass 1.37 GeV has very small πK-branching ratio, so its contribution is smaller then contributions of more massive resonances. The general tendency is that the contributions decrease with increasing the mass. We conclude that the

low energy approximation is reasonable and probably cannot be destroyed by inclusion of high energy effects. Using above estimates, we get

$$a_0^{3/2} - a_0^{1/2} = (0.22 \pm 0.02)\mu^{-1} \tag{5.2}$$

The small error illustrates consistency of the approach. It does not take into account, however, special reference to the FFGS model of the pion form factor, high-energy contributions to the unitarity condition, and higher partial waves ($l = 3, 5, \ldots$) in the t-channel backward scattering amplitude. The chiral theory in tree approximation gives a close prediction [15] $a_0^{3/2} - a_0^{1/2} = 0.21\mu^{-1}$. The estimates [17] based on the microscopic quark-gluon low energy dynamics of pseudoscalar mesons give $a_0^{3/2} - a_0^{1/2} = (0.33 - 0.60)\mu^{-1}$. The estimate (5.2) is in good agreement with the experimental data [18] $a_0^{3/2} - a_0^{1/2} = (0.14 - 0.37)\mu^{-1}$.

The kaon form factor evaluated according to eq.(4.1) is shown in Fig. 1 (curve C). The distinction from the naive estimates (A,B) is not very significant. The experimental data for the K^+ and K^0 electromagnetic form factors exist in the timelike region [2] and, for the charged kaon form factor, in narrow spacelike region at $t = 0$ [3]. In the timelike region, the kaon form factors are dominated by their isoscalar components, since in the resonance region typical values of the form factors, due to the vector-meson-dominance model, are $F(t) = O(m/\Gamma)$ where m and Γ are the vector meson masses and widths. In case of the ρ-meson (isovector contribution) $F_v(m_\rho^2) \approx 10$, whereas for isoscalar form factor the small ω-meson width implies $F_s(m_\omega^2) \approx 100$. The experimental data on the pion (isovector) form factor [1] and the charged kaon form factor [2] confirm these estimates. The isoscalar dominance of kaon form factors poses difficulties in extracting experimental information about isovector kaon form factor in the timelike region even if relative phases of isovector and isoscalar kaon form factors are known. In the spacelike region, the vector meson widths are not so important and one can expect $F_v(t) \approx F_s(t)$. Following this assumption, we compare our results with the experimental data [3] in Fig. 2. The agreement is rather good.

VII. CONCLUDING REMARKS

We constructed a one-dimensional integral representation for isovector kaon form factor in terms of the pion form factor and the backward πK^- scattering amplitude and derived, using the normalization condition at zero momentum transfer, a sum rule for the scattering amplitude. The estimate of the scattering lengths $a_0^{3/2} - a_0^{1/2} = 0.22\ \mu^{-1}$ is in agreement with the experimental data. The representation (4.1) satisfies low-energy unitarity, analiticity, and crossing symmetry.

The isovector kaon form factor numerically is very close to half of the pion form factor. The accurate calculations turn out to be close to estimate $F_K^v(t) = F_\pi(t)/2$ obtained assuming the ratio between the πK^- scattering amplitude and the pion form factor is a constant. This hypothesis is equivalent to the ρ-meson dominance in the πK^- scattering amplitude. It is really justified, since contributions to the sum rule of the lowest partial waves nearly cancel each other, so the right side of eq.(4.1) is determined mainly by the scattering lengths. This is a reason according to which the error in estimate of the scattering lengths is rather small. The accuracy of relation $F_K^v(t) = F_\pi(t)/2$ according to eq.(4.1) is $O(\Gamma/m)$. Deviations of the kaon form factor (4.1) from half of the pion form factor are smaller then 10% in the region of the ρ-meson peak, and negligible at $t < 0$.

The kaon form factor (curve C) decreases slightly faster as $t \to -\infty$ in comparison with the pion form factor. This is not reasonable, since strange quarks are more massive, and quarks inside of kaons are at smaller average distances (form factors of pointlike particles do not depend on momentum transfer). The numerical distinction, however, does not exceed one-two percents in the spacelike region, and is beyond accuracy of our approach. The second curve B in Fig.1 satisfies in fact that requirement. It was obtained assuming dominance of the integral term in the integral representation. We have found, however, by doing accurate estimates, that it is not the case.

The one-dimensional integral representation simplifies analysis of the kaon form factor behavior. In particular, it becomes apparent that asymptotics of the kaon form factor is determined by and the same as asymptotics of the pion form factor. In the region $|t| < 10^6$ GeV^2 the kaon form factor behavior is in agreement with the quark counting rules.

The integral representation (4.1) is based on the FFGS model for the pion form factor. It is clear, however, that specific feature of the model we used is eq.(3.3) only, so our results can be extended to any other model where similar equation exists.

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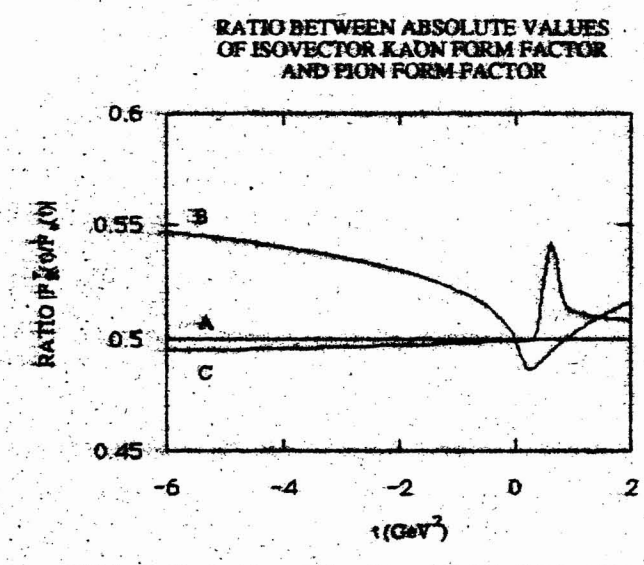

FIG.1. Ratio between absolute values of isovector kaon form factor and pion form factor: A - according to eq.(3.12), B - eq.(4.3), and C - eq.(4.1), assuming dominance of three lowest partial amplitudes in the integral term.

REFERENCES

1. C.Bebek *et al.*, Phys.Rev. **D17** (1978) 1693;
   B.Esposito *et al.*, Lett.Nuovo Cim. **28** (1980) 337;
   Phys.Lett. **67B** (1981) 239.
2. B.Esposito *et al.*, Nucl.Phys. **28** (1980) 337;
   B.Delcourt *et al.*, Phys.Lett. **99B** (1980) 257;
   F.Mane *et al.*, Phys.Lett. **99B** (1981) 261;
   P.M.Ivanov *et al.*, Phys.Lett. **107B** (1981) 297;
3. S.R.Amendolia *et al.*, Phys.Lett. **178B** (1986) 435.
4. R.G.Arnold *et al.*, Phys.Rev.Lett. **57** (1986) 174;
   R.C.Walker *et al.*, Phys.Lett. **B 224** (1989) 353; **240** (1990) 522;
   P.E.Bosted *et al.*, Phys.Rev.Lett. **68** (1992) 3841;
   G.Bardin *et al.*, Phys.Lett. **B 235** (1990) 123.
5. B.V.Geshkenbein and M.V.Terentyev, Yad.Fiz. **40** (1984) 758.
6. S.Dubnicka, Frascati report LNF-87/72, Frascati (Roma), 1987.
7. F.Iachello, A.Jackson, and A.Lande, Phys.Lett. **43 B** (1973) 191;
   G.Hohler *et al.*, Nucl.Phys. **114** (1976) 505;
   M.Gari and W.Krumpelmann, Z.Phys. **A 322** (1985) 689;
   S.Dubnicka, Nuovo Cim. **103A** (1990) 1417;
   P.L.Chung and F.Coester, Phys.Rev. **D 44** (1991) 229;
   P.Kroll, M.Schurmann, and W.Schweiger, Z.Phys. **A 338** (1991) 339.
8. W.Fraser and J.Fulco, Phys.Rev.Lett. **2** (1959) 365; Phys.Rev.**115** (1959) 1763; **117** (1960) 1609.
9. G.Gounaris and J.J.Sakurai, Phys.Rev.Lett. **21** (1968) 244.
10. P.S.Isaev and M.Sewerinski, Nucl.Phys. **22** (1961) 663;
    27 (1961) 148;
    P.S.Isaev, "Physics of Elementary Particles and Atomic Nuclei", **4** (1973) 731.
11. S.D.Drell and F.Zachariasen, "Electromagnetic Structure of Nucleons", 1961 (Oxford University Press).
12. D.V.Shirkov, V.V.Serebryakov, V.A.Mescheryakov, "Dispertion Theories of Strong Interactions at Low Energies", 1967 (Moscow, "Nauka").
13. A.V.Efremov, V.A.Mesheryakov and D.V.Shirkov, JETP **39** (1960)438.
14. W.Hogland *et al.*, Nucl.Phys. **B126** (1977) 109;

P.Estabrooks *et al.*, Nucl.Phys. **B133** (1978) 490;
D.Aston *et al.*, Nucl.Phys. **B296** (1988) 493.
15. M.K.Volkov and V.N.Pervushin, "Nonlinear Quantum Theories, Dynamical Symmetries, and Physics of Mesons", 1978 (Moscow, "Atomizdat").
16. J.Gasser and H.Leutwyler, Ann.Nucl.Phys. **158** (1984) 142.
17. J.Weinstein, Phys.Rev. **D47** (1993) 911.
18. N.O.Johannesson and J.Petersen, Nucl.Phys. **B68** (1973) 397;
M.J.Matison *et al.*, Phys.Rev. **D9** (1974) 1872;
C.B.Lang *et al.*, Nuovo Cimento **A41** (1977) 73;
A.Karabouraris and G.Shaw, J.Phys. **G6** (1980) 583;
O.Drumbajs *et al.*, Nucl.Phys. **B216** (1983) 277.

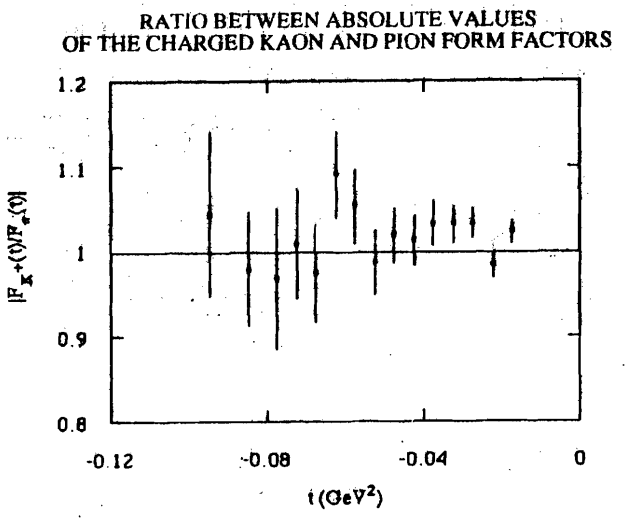

Fig. 2. Ratio between charged kaon and pion form factors assuming that isovector and isoscalar kaon form factors are equal at small t. The curve shown here is the same as the curve C in Fig. 1. The experimental data are from ref. [3].

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