ANTIKAON-NUCLEON SCATTERING AT LOW ENERGIES

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A separable, quasirelativistic, potential model is used to describe $\bar{K}$N elastic and inelastic scattering. Free parameters of the model are set to reproduce the scattering data up to 300 MeV/c momentum and the virtual $Y_0(1405)$ bound state energy and width. Various inelastic processes occurring below the $\bar{K}$N threshold and observed in the kaon nuclear interactions are studied. The dependence on the phenomenological parameters is discussed.

1. *Introduction*

The stimulus to study the negative kaon interactions at low energies, apart from the interest of elementary particle physics, comes from nuclear physics. It has been argued since a long time that kaon is the best tool to study the nuclear surface composition [1]. This requires a knowledge of kaon nucleon multichannel scattering amplitudes off the energy shell. The aim of this paper is to provide a simple model of such amplitudes. Here, the model is tested on the available two body scattering data. A further aim is to use it for a description of few body kaonic processes and for nuclear interactions of the kaons.

Our analysis is done in terms of a quasi relativistic Schrödinger equation. The interaction is introduced in the form of separable multichannel potential. Some properties of the equation are given in Sec. 2. In Sec. 3 we discuss the procedure and numerical results which consists of a reproduction of the kaon-nucleon scattering data and extrapolation of the amplitudes below the $\bar{K}$N threshold. A special attention is paid to the behaviour of the resonant state. A dependence of solution on the parameters of the potential is also studied and a fairly strong sensitivity of the off shell behaviour to the range of the forces is found.

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2. The model

The two main features of K^-N reactions at low energies are a short range of the forces and the presence of open hyperon pion channels. The forces are believed to be due mainly to the exchange of vector mesons, thus being of the range $1/\mu \approx 0.25$ fm, [2]. However, a longer range component arising from the two pion exchange is also possible. This is left open and the range parameters are free parameters which eventually turn out to be of intermediate magnitude. This is supported by a number of phenomenological analyses, [2–7], which show that the S wave scattering dominates the picture up to 300–400 MeV/c of the kaon laboratory momentum. We restrict ourselves to this wave only. The coupling of the initial K^-p channel to another kaonic K^0n and the hyperonic $\Sigma\pi$, $\Lambda\pi$ channels is pretty strong giving rise to a resonant state of the $\Sigma\pi$ pair. We shall consequently consider only the S wave in all the channels. The three body $\Lambda\pi\pi$ channel is also open but this is known to be of no importance in the region of interest, being usually disregarded.

Within the two body space the plane wave states will be normalized in a Lorentz invariant way

$$\langle i\vec{p}_1\vec{p}_2|j\vec{p}_1'\vec{p}_2'\rangle = \frac{(2\pi)^2}{\sqrt{S}} E_i(p_1)\omega_i(p_2)\delta(\vec{p}_1-\vec{p}_1')\delta(\vec{p}_2-\vec{p}_2')\delta_{ij}$$
$$= \delta(\vec{P}-\vec{P}')\,\{N_i^2(q)\delta(\vec{q}-\vec{q}')\}\delta_{ij}, \tag{1}$$

where i, j are channel indices and P, q are the total and relative momenta.

Here from we use the centre of mass system and thus the momentum conservation delta factorizes out in all equations. The normalization factor becomes

$$N_i^2 = (2\pi)^2 E_i\omega_i/(E_i+\omega_i),$$
$$E_i = (m_i^2+q^2)^{1/2}, \qquad \omega_i = (\mu_i^2+q^2)^{1/2}. \tag{1a}$$

The scattering matrix will be determined from the Lippman Schwinger equation. If we denote

$$\langle\vec{p}_1\vec{p}_2|T(E)|\vec{p}_1'\vec{p}_2'\rangle = T(\vec{q},\vec{q}\,';E)\delta(\vec{P}-\vec{P}'),$$
$$\langle\vec{p}_1\vec{p}_2|V|\vec{p}_1'\vec{p}_2'\rangle = V(\vec{q},\vec{q}\,')\delta(\vec{P}-\vec{P}'),$$

this becomes

$$T_{ij}(\vec{q},\vec{q}\,';E) = V_{ij}(\vec{q},\vec{q}\,')+\sum_k \int \frac{d_3\vec{q}\,''}{N_k^2(q'')}\,\frac{d_3\vec{q}\,'''}{N_k^2(q''')}$$
$$\times V_{ik}(\vec{q},\vec{q}\,'')G_k(\vec{q}\,'',\vec{q}\,''';E)T_{kj}(\vec{q}\,''',\vec{q}\,';E). \tag{2}$$

We need a quasi-relativistic approach because of pions in the inelastic channel. A number of relevant methods may be found in the literature, differing by a choice of the Green function G in Eq. (2). These are various reductions of the Bethe-Salpeter equations to the equal time formalism. For some the probabilistic interpretation seems to be unclear although relativistic invariance is achieved.

We would rather prefer to loose partly the latter so we project on the positive energy solutions only. The Green function is then normalized by the propagation condition

$$\lim_{t\to t_0} \hat{G}(t, t_0) = \hat{1}/i = \int \frac{dE}{2\pi} \hat{G}(E) = \int \frac{dE}{2\pi} \cdot \sum_n \varphi_n \bar{\varphi}_n (E-E_n)^{-1}, \tag{3}$$

which is equivalent to the completeness of the set of positive frequency solutions φ_n. With the normalization (1) we obtain

$$\langle iq' | \hat{G}(E) | \vec{q} j \rangle = \delta_{ij} N_i^2(q) \delta(\vec{q}-\vec{q}\,') \left[E - E_i(q) - \omega_i(q) + i\varepsilon\right]^{-1}, \tag{4}$$

and the integral equation simplifies to the form

$$T_{ij}(\vec{q}_i, \vec{q}_j; E) = V_{ij}(\vec{q}_i, \vec{q}_j) + \frac{1}{(2\pi)^2} \sum_k \int \frac{d_3 \vec{q}_k}{E_k(q_k)\omega_k(q_k)}$$

$$\times V_{ik}(\vec{q}_i, \vec{q}_k) \frac{1}{[E - E_k(q_k) - \omega_k(q_k) + i\varepsilon]} T_{kj}(\vec{q}_k, \vec{q}_j; E). \tag{5}$$

The normalization (1) was chosen so that the residue in the Green function (4), when integrated over momenta, equals unity. Hence the asymptotic behaviour of the wave function is

$$\Psi_f = \Phi_i + G_f T_{fi} \Phi_i \approx e^{i\vec{q}_i \vec{x}} - T_{fi}(q_f, q_i; E) \frac{e^{iq_f r}}{r},$$

and the Lorentz invariant expression for the current yields the cross section

$$\frac{d\sigma_{fi}}{d\Omega} = \frac{q_f}{q_i} \cdot T_{fi}(q_f, q_i; E)|^2. \tag{6}$$

3. *Results*

Experiments distinguish the initial K^-p channel coupled to the other kaonic state $\bar{K}^0$n. A suitable combination forms the pure isospin $I = 0, 1, I_3 = 0$ states. These are coupled to $\Sigma\pi$ $I = 0, 1$ and $\Lambda\pi$ $I = 1$ states, which together form a five dimensional basis in the isospace: $(K\mathrm{N})_0, (K\mathrm{N})_1, (\Sigma\pi)_0, (\Sigma\pi)_1, (\Lambda\pi)_1$. The other I_3 components for example K^-n states are not accessible in the two body experiments but appear in the K^- deuteron or K^- nucleus scattering. The potential V is assumed to be an isospin invariant 2×2 for $I = 0$ and 3×3 for $I = 1$ matrix. This symmetry is broken, however, by Coulomb interactions and by the mass splitting of the isospin multiplets. These will be introduced into the Green functions. In the region of interest far from the $\Sigma\pi$ threshold the mass splitting of the Σ and π multiplets is not relevant but the 4.9 MeV difference of K^-p and K^0n thresholds should be considered.

The Green function is thus diagonal in the channel representation (K^-p), (K^0n), and $(\Sigma\pi)_{1,0}$, $(\Lambda\pi)_1$ while the potential is invariant in the isospin representation

$$(KN)_0 = 1/\sqrt{2}\,(K^-p) - 1/\sqrt{2}\,(K^0n),$$

$$(KN)_1 = 1/\sqrt{2}\,(K^-p) + 1/\sqrt{2}\,(K^0n), \tag{7}$$

of the kaonic component. The corresponding transformation is achieved by a matrix U

$$G_{KL} = \sum U_{Ki} G_i U^+_{Li}, \qquad V_{ij} = \sum U^+_{Ki} V_{KL} U_{Lj}, \tag{8}$$

where capital indices denote isospin representation. In the kaonic subspace $U = 1/\sqrt{2}\begin{pmatrix}1 & -1\\ 1 & 1\end{pmatrix}$ follows from Eq. (7). In the hyperon pion channels the isospin representation is used throughout.

We use potential with one separable term only

$$V^I_{KL}(\vec{q}, \vec{q}') = \mathscr{V}_K(q)\lambda_{KL}(I)\mathscr{V}_L(q'), \tag{9}$$

with Yamaguchi formfactors

$$\mathscr{V}_K(q) = \frac{\kappa_K^2}{\kappa_K^2 + q^2}, \tag{10}$$

and this corresponds to the S wave only. The potential matrix $\lambda(I)$ is real and symmetric to guarantee hermiticity and time reversal invariance. For simplicity reasons the inverse range parameters κ are chosen to depend on the channel (and not isospin) index only. The free parameters are then: two inverse ranges α_K, α_Σ and three depth parameters $\lambda_{KL}(0)$ for $I = 0$ and six $\lambda_{KL}(1)$ plus one κ_Λ for $I = 1$ that is twelve parameters altogether.

In the channel representation the solution of Eq. (5) has the form:

$$T_{ij}(q, q'; E) = \mathscr{V}_i(q) t_{ij}(E) \mathscr{V}_j(q'), \tag{11}$$

where

$$[t^{-1}(E)]_{ij} = [\lambda^{-1}]_{ij} - R_i(E)\delta_{ij},$$

$$R_i(E) = \frac{1}{\pi}\int_0^\infty \frac{E}{E_i(q)\omega_i(q)} \frac{\mathscr{V}_i^2(q) q^2 dq}{E - E_i(q) - \omega_i(q) + i\varepsilon}.$$

Since the masses of barions m_i are much larger than those of mesons μ_i, we may approximate the functions R_i by the following expressions

$$R_i(E) \approx \frac{1}{\pi}\int_0^\infty \frac{E}{m_i\omega_i(q)} \frac{\mathscr{V}_i^2(q) q^2 dq}{E - m_i - q^2/2m_i - \omega_i(q) + i\varepsilon}. \tag{12}$$

For the Yamaguchi formfactors this integral may be performed explicitly. To do it let us introduce abbreviations:

$$\tau^{(i)}_{1,2}(E) = \frac{1}{\mu_i}[-m_i \pm (2Em_i + \mu_i^2 - m_i^2)^{1/2}],$$

$$k^i_{1,2} \equiv k^i_{1,2}(E) = \mu_i\{[\tau^{(i)}_{1,2}(E)]^2 - 1\}^{1/2},$$

$$\alpha^i_1(E) = \ln\left[\tau^{(i)}_1(E) + \frac{1}{\mu_i} k^i_1(E)\right],$$

$$\alpha^i_2(E) = \ln\left|\tau^{(i)}_2(E) + \frac{1}{\mu_i} k^i_2(E)\right|,$$

$$\gamma_i = \ln\left[\frac{\kappa_i - (\kappa_i^2 - \mu_i^2)^{1/2}}{\kappa_i + (\kappa_i^2 - \mu_i^2)^{1/2}}\right],$$

$$M^i_{1,2}(E) = -\frac{1}{\pi} k^i_{1,2}(E)\alpha^{(i)}_{1,2}(E)\mathscr{V}^2_i(k^i_{1,2})$$

$$-\frac{\mu_i^2}{4(\kappa_i^2 - \mu_i^2)}\mathscr{V}_i(k^i_{1,2})\left\{\frac{\kappa_i^2}{\pi\mu_i}\tau^{(i)}_{1,2}(E)\left[2 + \frac{\gamma_i \mathscr{V}_i(k^i_{1,2})}{\kappa_i^3(\kappa_i^2 - \mu_i^2)^{1/2}}[2\kappa_i^4 + (k^i_{1,2})^2\mu_i^2 - \kappa_i^2\mu_i^2]\right]\right.$$

$$\left. + \frac{1}{\mu_i^2\kappa_i}\mathscr{V}_i(k^i_{1,2})(\kappa_i^2 - \mu_i^2)[\kappa_i^2 - (k^i_{1,2})^2]\right\}.$$

The function $R_i(E)$ is then

$$R_i(E) = -ik^i_1(E)\frac{E}{(2m_iE + \mu_i^2 - m_i^2)^{1/2}}\mathscr{V}^2_i(k^i_1) + P_i(E), \tag{13}$$

where

$$P_i(E) = -\frac{E}{\pi(2m_iE + \mu_i^2 - m_i^2)^{1/2}}[M^i_1(E) - M^i_2(E)].$$

Let us comment on the effect of the Coulomb forces. To calculate it exactly the procedure of separable model must be repeated with the free Green function G_0, (3), replaced in Eq. (2) by the Coulomb Green function. This is performed more easily in the coordinate representation. For the region of energy under study the Coulomb force is significant in the K^-p channel and close to the threshold only. Thus we use nonrelativistic form of the Coulomb corrections. The exact solution of this problem within our Eq. (5) is also simple but somewhat lengthy. As the details will be published elsewhere we give the result which consists of a change of Eq. (13):

$$R^{\text{cul}}(E) = R(E) + 2\mu\alpha\left\{-\ln\left(\frac{-4ik_1}{\kappa}\right) + i\gamma\sum_{n=1}^{\infty}\frac{1}{n(-i\gamma + n)}\right\}, \tag{14}$$

where $\gamma = \alpha/\sqrt{2\mu E}$. Only the leading term in $\mu\alpha/\kappa$ is retained. This formula is valid on the whole cut energy plane and yields also energies and widths of the atomic quasibound states of the K^- and p. These are given by positions of the poles in the elastic K^-p amplitude, that is, by the equation

$$\frac{1}{n-\gamma(E)} = \frac{1}{2\mu\alpha t_{K^-p}(E)} + \ln\left(\frac{-4ik_1(E)}{\kappa}\right) - i\gamma(E)\sum_{m\neq n}\frac{1}{m(m-i\gamma(E))}, \tag{15}$$

where n is the principal atomic quantum number. Keeping the first dominant term only a well known formula is obtained

$$E_n = \frac{\mu\alpha^2}{2n^2} + \frac{2\mu^2\alpha^3}{n^3}\, t(E_n), \tag{16}$$

which gives the strong shift to 15% accuracy. The relativistic corrections to the lowest level are still 10^{-2} of the strong shift. The other change to be made concerns the relation of the matrix (10) to the experiment. The on shell scattering matrix which enters formula (6) for the cross section is now

$$T_{fi} = \mathscr{V}_f(k_1^f(E))t_{fi}(E)\mathscr{V}_i(k_1^i(E))\Gamma(1+\gamma_f)e^{\frac{\pi}{2}\gamma f}$$
$$\times\Gamma(1-i\gamma_i)e^{\frac{\pi}{2}\gamma i} + \delta_{if} f_{\text{Coulomb}}(\theta) \tag{17}$$

which together with Eq. (14) are a slight generalization of the well known effective range formulas.

Let us now turn to the description of the experimental data. To determine the parameters of this model the following conditions will be met:

a) reproduction of the energy and width of the Y_0 resonance,

b) best fit to the low energy K^-p elastic and inelastic data.

The requirement a) leads to the conditions

$$\text{Re}\; T^{-1}_{(\Sigma\pi)_0(\Sigma\pi)_0}(E_R) = 0,$$
$$\frac{d}{dE}\,\text{Im}\; T^{-1}_{(\Sigma\pi)_0(\Sigma\pi)_0}(E_R) = \Gamma/2, \tag{18}$$

with E_R = 1406 MeV and I = 42 MeV chosen. These turn out to yield strong restrictions on κ_{KN}, $\kappa_{\Sigma\pi}$ and $\hat{\lambda}(0)$. The best fit to the scattering data was made in a simplified way by comparing our reaction matrix

$$K_{ij} = \mathscr{V}_i(q)k_{ij}(E)\mathscr{V}_j(q),$$
$$k_{ij}^{-1} = \lambda_{ij}^{-1} - \delta_{ij}P_i, \tag{19}$$

with phenomenological sets of the reaction matrices found by various authors [3–7]. This may be done at a given, fixed, energy as our K matrix is energy dependent in a complicated way while the phenomenological matrices are given either in the scattering length approximation or in the effective range approximation

$$K_{ij}^{-1}(E) = K_{ij}^{-1}(0) - \tfrac{1}{2} r_i q_i^2 \delta_{ij}.$$

The difference of the models manifests itself mainly through an energy dependence of the nondiagonal terms which appear in our formalism. If we keep the resonance conditions (18) fixed it becomes difficult to obtain a reasonable fit to those solutions for which det $|K_{I=0}(0)| < 0$ [3, 5, 7, solution "b"]. This happens so because the sign of the determinant introduces correlation of the relative signs of the scattering lengths and the effective ranges. This phenomenon known already in the effective range approach [8] occurs in this model too.

TABLE I

Potential matrix parameters of this paper and the reaction matrices of Ref. [7], in fermis

$$\begin{matrix} & (KN)_0 & (\Sigma\pi)_0 \end{matrix}$$

$$K(I=0) = \begin{pmatrix} -2.27, & -1.14 \\ , & -0.90 \end{pmatrix}; \ \lambda(I=0) = \begin{pmatrix} -0.48, & -0.42 \\ , & 0.25 \end{pmatrix}$$

$$\begin{matrix} & (KN)_1 & (\Sigma\pi)_1 & (\Lambda\pi)_1 \end{matrix}$$

$$K(I=1) = \begin{pmatrix} 0.183 & -0.649 & -0.476 \\ & -0.324 & -0.141 \\ & & 0.550 \end{pmatrix}$$

$$\lambda)I=1) = \begin{pmatrix} 1.42 & -2.09 & -1.73 \\ & 2.03 & 1.99 \\ & & 1.07 \end{pmatrix}$$

We find good overall agreement with the solutions of A. D. Martin and G. G. Ross [4] and B. R. Martin solution *a*, [7][1]. The parameters of the best fit are given in the Table I together with the B. R. Martin reaction matrices which were used for this fit. The reciprocal range parameters are $\kappa_{KN} = 1094$, $\kappa_{\Sigma\pi} = 799$ and $\kappa_{\Lambda\pi} = 700$ MeV and the ranges of corresponding potentials are roughly $2/\kappa_l$. The solution is pretty stable with respect to the energy at which comparison is made up to $E = 1480$ MeV. The actual fit was made at $E = 1470$ MeV or $q_{KN} = 240$ MeV/c. For higher energies of comparison other, unphysical, solutions with comparable χ^2 arise too. Various scattering matrices $T(E)$ on the mass shell, but without the Coulomb penetration factors of Eq. (17) and corrections (14), are drawn in Fig 1. In Figs 1a, b the elastic amplitudes K^-p and $(\Sigma\pi)_0$ coupled strongly to the resonant $I = 0$ amplitude are given. The continuous curves come from the best fit parameters. Those dashed correspond to the reverse ranges κ_{KN} and $\kappa_{\Sigma\pi}$ reduced by 20% and the unchanged strength parameters λ. The amplitudes in the scattering region, above the threshold, are very stable but the effect on the resonance is significant. This sensitivity to the off shell behaviour is an unpleasant factor for the many body low energy kaon interactions. The lesson is that the amplitudes below the threshold must be treated with care and possibly checked in an independent way.

The elastic $I = 1$ amplitudes shown in Figs 1c, d are almost κ independent.

[1] The authors are grateful to Dr. B. R. Martin for sending us his results prior to publication.

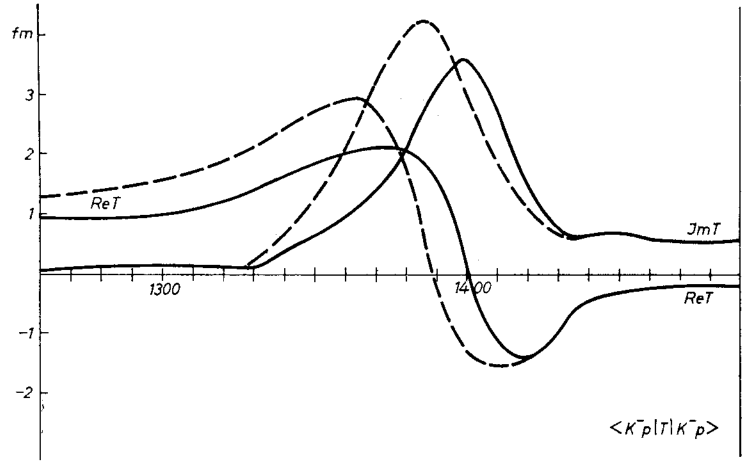

Fig. 1a. The diagonal matrix element $T_{(K^-p),\,(K^-p)}$ versus energy. Best fit parameters. The dashed lines represent the same element but with κ reduced by 20%

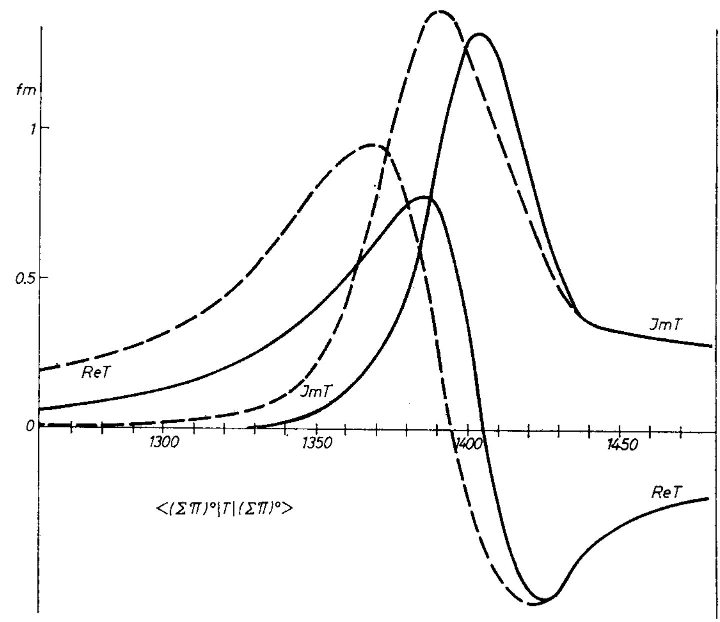

Fig. 1b. The diagonal matrix element $T_{(\Sigma\pi)_0,\,(\Sigma\pi)_0}$ versus energy. Best fit parameters. The dashed lines represent the same element but with κ reduced by 20%

Finally various cross sections and branching ratios are given in Figs. 2–7, with the corresponding experimental data. Various characteristics of the sigma hiperon production are given in Figs 2, 4, 5, 8. The sensitivity of amplitudes to the range of forces in the subthreshold region is shown in Figs 2 and 8 where the branching ratios:

$$R_{+/-} = \text{Yield } (\Sigma^+\pi^-)/\text{Yield } (\Sigma^-\pi^+)$$
$$R_{pn} = \text{Yield } (\Sigma^+\pi^- + \Sigma^-\pi^+)/\text{Yield } (\Sigma^-\pi^0)$$

are given.

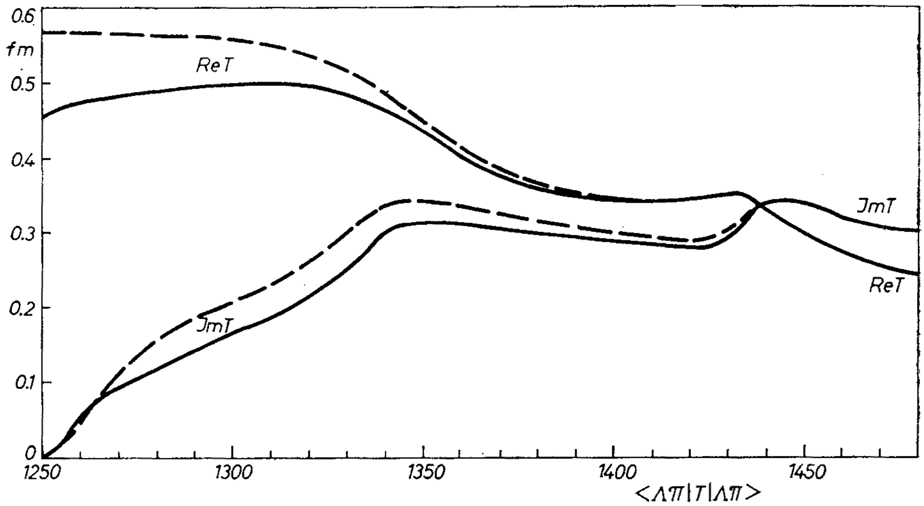

Fig. 1c. The diagonal matrix element $T_{(\Lambda\pi)_1\,(\Lambda\pi)_1}$ versus energy (see caption to Fig. 1a)

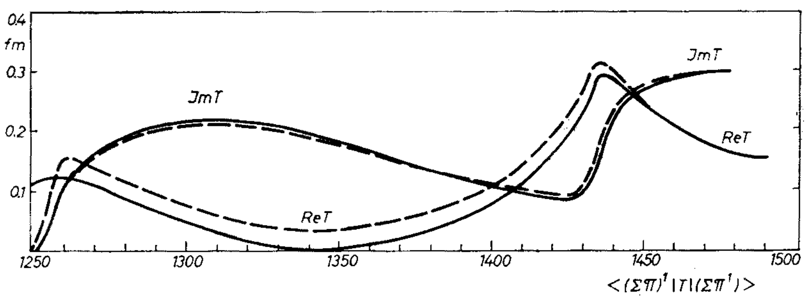

Fig. 1d. The diagonal matrix element $T_{(\Sigma\pi)_1\,(\Sigma\pi)_1}$ versus energy (see caption to Fig. 1a)

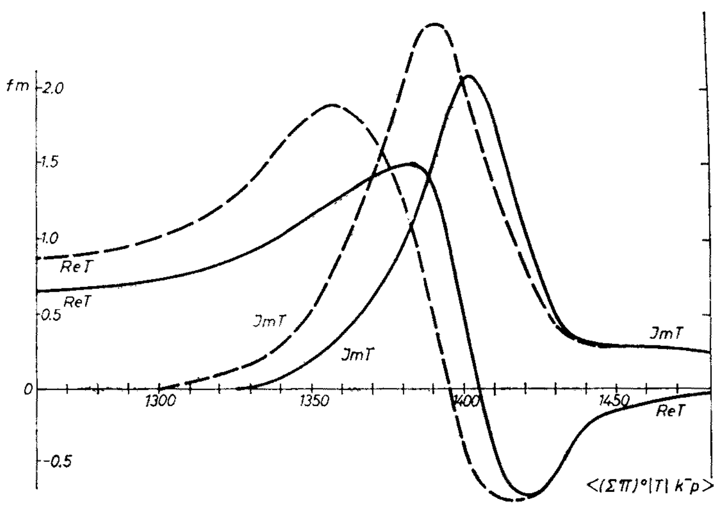

Fig. 2. The non diagonal matrix element $T_{(\Sigma\pi)_0\,(K^-p)_0}$ versus energy, in the isospin 0 state (see caption to Fig. 1).

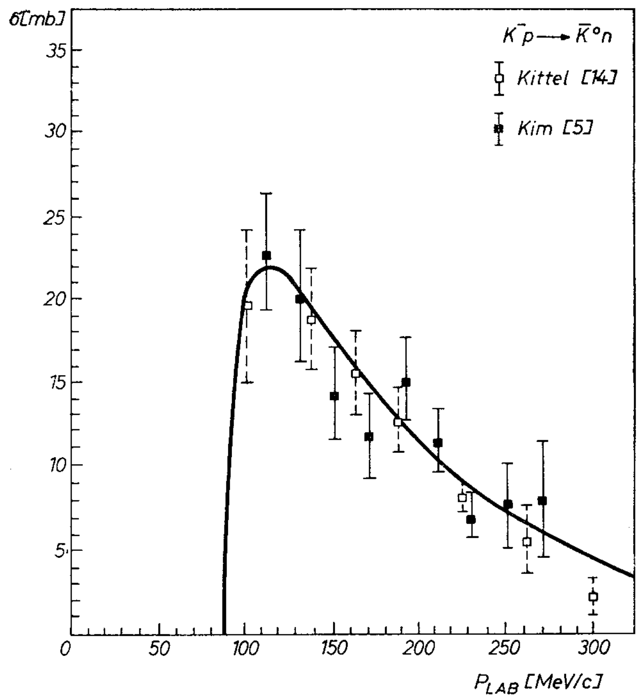

Fig. 3. Calculated cross section for $K^-p \rightarrow \overline{K}^0n$ versus kaon laboratory momentum

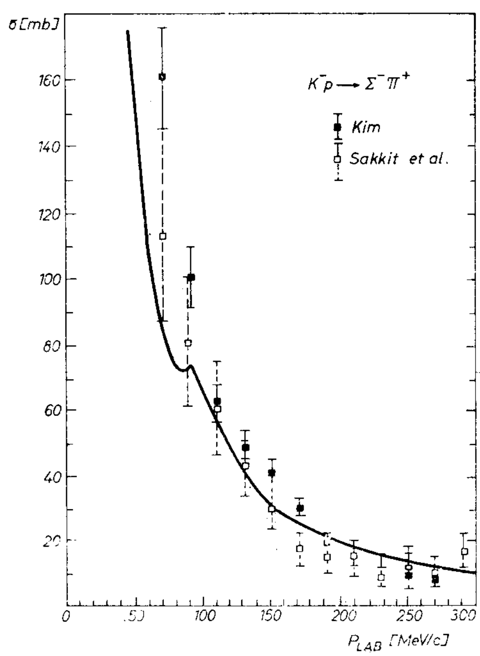

Fig. 4. Calculated cross section for $K^-p \rightarrow \Sigma^-\pi^+$ versus kaon laboratory momentum

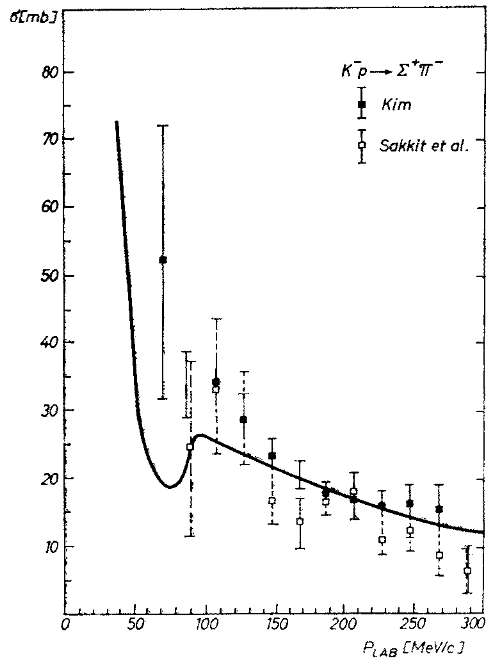

Fig. 5. Calculated cross section for $K^{-}p \rightarrow \Sigma^{+}\pi^{-}$ versus kaon laboratory momentum

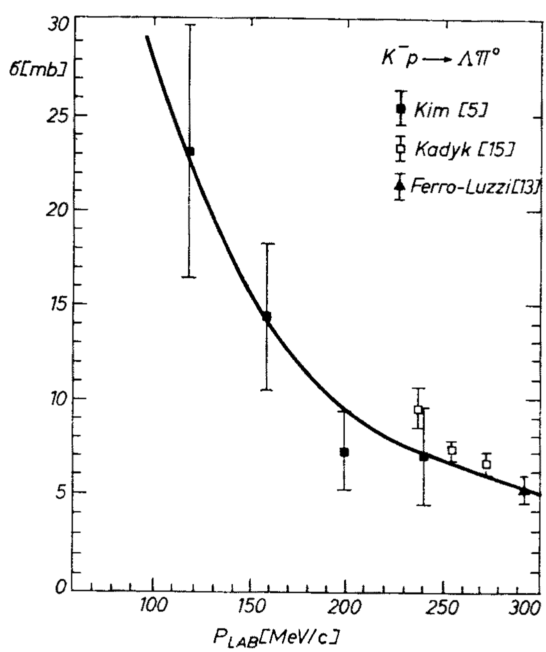

Fig. 6. Calculated cross section for $K^{-}p \rightarrow \Lambda\pi^{0}$ versus kaon laboratory momentum

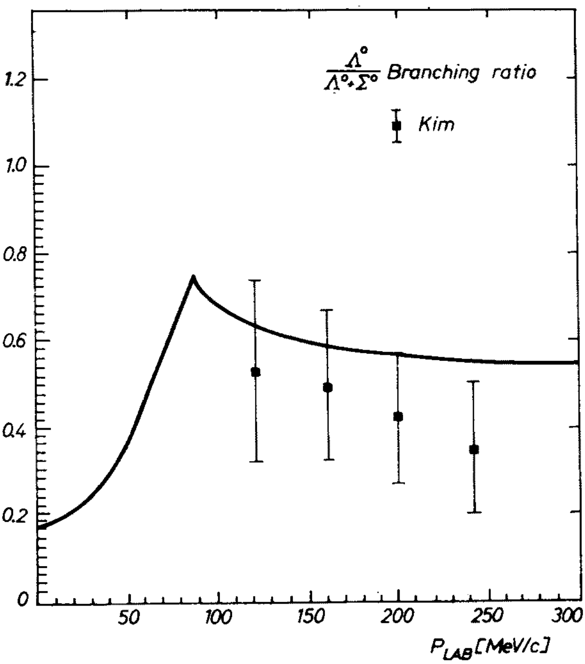

Fig. 7. Calculated branching ratio $(K^-p \to \Lambda\pi^0)/[(K^-p \to \Lambda\pi^0)+(K^-p \to \Sigma^0\pi^0)]$ versus kaon laboratory momentum

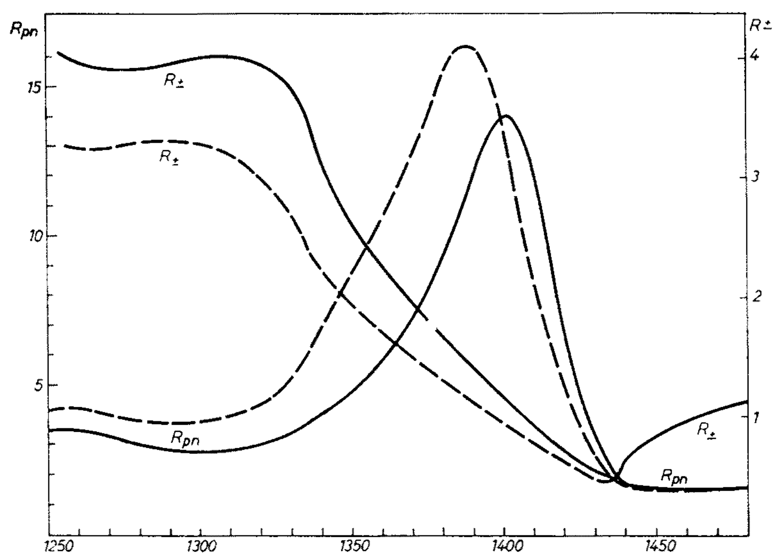

Fig. 8. Branching ratios R_{+-} and R_{pn}. See caption to Fig. 1

These ratios have been used to characterize nuclear kaonic interactions and are a key to study the neutron skin effect in nuclei [9, 10]. The dependence on κ is moderate and probably cannot be used to explain the high experimental number for $R_{pn} \approx 20$ in light emulsion nuclei, [9]. A plausible explanation as noted in Ref. [11] lies in a poor knowledge of $I = 1$, $\bar{K}N \rightarrow \Sigma\pi$ amplitude at low (zero) energies. Many body effects are also not negligible in the nuclear case, [12]. The Coulomb effects given as corrections in (15) and multiplicative penetration factor in formula (17) are important only in the very low energy region and in the atomic region. For relative momenta $q_{KN} \approx 100$ MeV or $q_{KN} \approx i \cdot 100$ MeV below threshold only the logarithm is nonnegligible. Around the resonance Y_0 it becomes of some significance as Re $t^{-1}_{\bar{K}N,\bar{K}N} \approx 0$ and causes $\sim$0.5 MeV upward shift of the 27 MeV resonance binding energy. This value of energy shift is accurate to about 25% as only the first term in the development of the Coulomb correction in (q/κ) was used. The accuracy of this development in thre atomic region is 0.5%.

Let us comment also on the stability of the model itself. It was based on the simple form (4) for the Green function. If we choose another, for example Klein Gordon equation's Green function for mesons, the contribution from negative frequency states appears. This, however, for the energies of interest contributes almost energy independent term to the real part of the function R of Eq. (12). The net result is a (small) renormalization of the coupling constant λ, only. Refinements of the potential in terms of new parameters do not seem necessary with the data at hand.

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