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ИНСТИТУТ ЯДЕРНОЙ ФИЗИКИ СО АН СССР

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THE (^{3}He, T) REACTION ON NUCLEI IN THE Δ-ISOBAR REGION

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ABSTRACT

The cross-section for (^{3}He, t) reaction on nuclei at high excitation energy is calculated in optical approximation for incoming and outgoing ions. For high momentum transfer $qR \gg 1$ the cross-section is expressed in terms of imaginary part of a pion propogator at nuclear surface. Comparison with the data demonstrates that the approximation is good only for comparatively low kinetic energy of ^{3}He.

1. INTRODUCTION

The spin-isospin response of a nucleus has been intensively studied in the (*p*, *n*) reaction with 200 MeV protons [1,2] (and ref. therein). Recently, the (^{3}He, t) reaction has been used to study the spin-isospin response of a nucleus at higher energy transfer where the internal degrees of freedom of a nucleon can be excited [3,4].

The data show clearly the difference between the objects excited from a single nucleon and from a nucleus. The difference is both in the peak positions and the heights of the peaks. According to conventional Glauber theory of quasifree scattering the height of the peak must be lower in the case of a carbon than in case of a single nucleon due to strong absorption of a projectile. The data, however, show the height of the peak on ^{12}C bigger than on a nucleon. The shift of the position cannot be explained as well by Fermi motion in quasifree mechanism.

This situation is similar to excitation of giant resonances at low energy where the transition strength is concentrated in a single collective state rather than in separate particle-hole states. The question now is—what kind of states can form the resonance at high energy. The obvious candidate is the Δ-isobar. This resonance gives a strong peak in the (^{3}He, t) reaction on a single nucleon. The collective state, however, can hardly be built as a superposition of pure delta-hole states. The reason is in very short lifetime of a Δ. During its life the Δ has no time to propagate further than to a neighbour nucleon, therefore, it does not feel the nuclear potential to form the state sensitive to the shape of a nucleus.

From the other hand, this large width of a delta can lead to formation of another kind of collective states. The delta decays mainly through pion emission. The pion due to large width of a delta is still in resonance and can be reabsorbed by another nucleon forming again a delta-isobar but in different point. In this way we obtain an excitation propagating through nuclear medium which is a superposition of the delta-hole and pionic degrees of freedom. At very high density, where deconfinment of quarks occurs, this excitation transforms into a longitudinal spin-wave in quark matter.

This picture is very close to that proposed in the pioneering paper of T.E.O. Ericson and J. Hüfner [5] for the scattering of real pions off a nucleus.

Following this line, in the next part we express the cross-section for the (^{3}He, t) reaction in terms of a pion propagator and in the third part we discuss shortly the properties of a pionic propogator inside a nucleus. The similar investigation has been done in [6] for the real pions. For virtual pions the Chew-Low model is not sufficient because we need an amplitude reproducing the nonresonant partial waves as well.

For the (^{3}He, t) reaction the similar approach has been developped in [7, 8] where nonrelativistic response function of a nucleus has been used.

2. THE CROSS-SECTION OF THE (^{3}He, t) REACTION ON A NUCLEUS

The data on $p(^3\mathrm{He}, t)\Delta^{++}$ reaction can be well described by one-pion exchange as well as $p(p, n)\Delta^{++}$ reaction. Good fit of the data can be obtained using the parameter of the formfactors of πNN and $\pi N\Delta$ vertexes $\Lambda = 630$ MeV.

Let us forget for the moment about distortion of the projectile-ejectile waves and start from the PWIA. The general expression for the cross-section can be written as follows

$$\frac{d^2\sigma}{d\Omega dE_t} = \frac{1}{16\pi^2}\frac{p_t}{p_{\mathrm{He}}}\sum_f \overline{|M_{0f}|^2}\,\delta(\omega + E_0 - E_f)\,d\varrho_f, \qquad (1)$$

where M_{0f} is the matrix element of the T-matrix. For one-pion exchange the matrix element is of the form

$$M_{0f} = -\sqrt{2}\,\frac{m_{\mathrm{He}}}{m}\,g_N(q^2)F(q^2)\bar{U}(p_t)\gamma_5 U(p_{\mathrm{He}})\,G_\pi^0(q^2)\cdot\Gamma_{\pi N\Delta}\,, \qquad (2)$$

here, the ^{3}He and t are treated as relativistic particles of spin 1/2. $F(q^2)$ is the formfactor of the (^{3}He, t) vertex. $G_\pi^0(q^2)$ is pionic propogator. The cross-section can be drawn graphically as in Fig. 1,

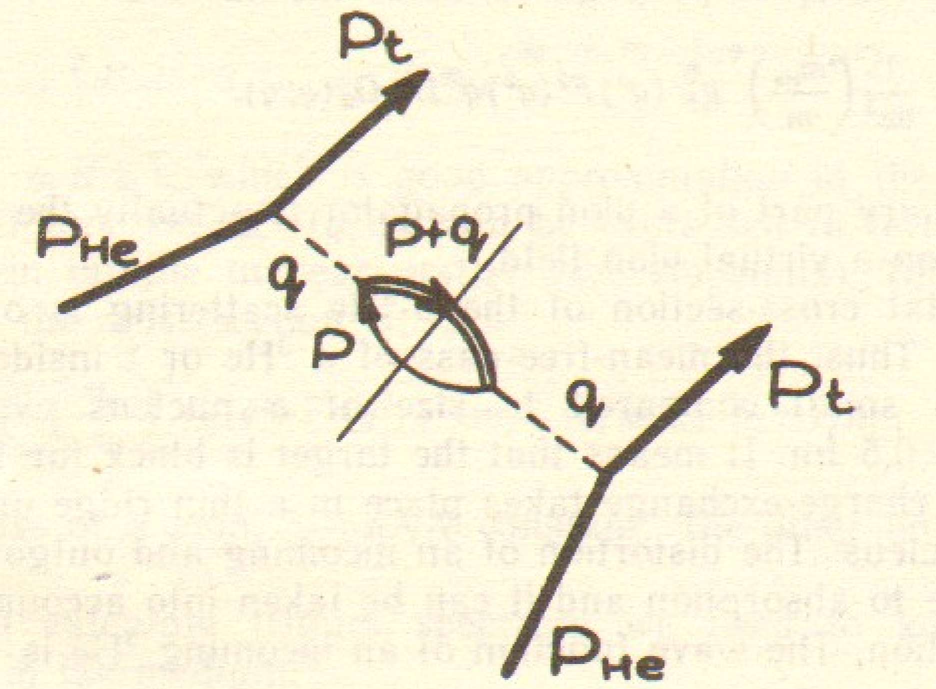

Fig.1. Matrix element squared for isobar production in nuclear matter.

with the expression for $\overline{|M_{0f}|^2}$ being

$$\overline{|M_{0f}|^2} = 2(-q^2)\left(\frac{m_{\mathrm{He}}}{m}\right)^2 g_N^2(q^2)F^2(q^2)\left(G_\pi^0(q^2)\right)^2\overline{|\Gamma_{\pi N\Delta}|^2}. \qquad (3)$$

The sum over final states in (1) includes the sum over the momentum of the hole $\vec{k}$ weighted with $\sum|\varphi_\nu(\vec{k})|^2$, where $\varphi_\nu(\vec{k})$ is the wave-function of the state occupied by a nucleon.

The matrix elements of $\pi N\Delta$-vertexes, together with the δ-function, form exactly the imaginary part of the pion self-energy inside a finite nucleus. Combining (1) and (3) we obtain the cross-section

$$\frac{d^2\sigma}{d\Omega dE_t} = \frac{1}{8\pi^3}\frac{p_t}{p_{\mathrm{He}}}\left(\frac{m_{\mathrm{He}}}{m}\right)^2 g_N^2(q^2)F^2(q^2)\,q^2\left(G_\pi^0(q^2)\right)^2 \mathrm{Im}\,\Pi(\vec{q},\vec{q})\,, \qquad (4)$$

$\Pi(\vec{q},\vec{q}) = \int d^3r\,d^3r' e^{-i\vec{q}(\vec{r}-\vec{r}')}\,\Pi(\vec{r},\vec{r}')$ is the self-energy of a pion inside a nucleus. For the infinite matter

$$\Pi(\vec{q},\vec{q}) = \frac{A}{n}\Pi(\vec{q}), \qquad (5)$$

where A is the number of nucleons and n is nuclear density.

Since the imaginary part of the bare pion propagator at given kinematics is equal to zero, we obtain the final expression for the cross-section by changing $G^0_\pi(g^2 \,\mathrm{Im}\,\Pi(\vec{q},\vec{q})\,G^0_\pi(q^2) \to \mathrm{Im}\,G_\pi(q^2)$, where $G_\pi(q^2)$ is the pion propagator inside the nucleus

$$\frac{d^2\sigma}{d\Omega dE_t} = \frac{1}{8\pi^3}\left(\frac{m_{He}}{m}\right)^2 g^2_N(q^2) F^2(q^2) q^2 \,\mathrm{Im}\, G_\pi(q,q). \qquad (6)$$

The imaginary part of a pion propagator is actually the response of a nucleus on a virtual pion field.

The total cross-section of the p-^{3}He scattering is of the order ~ 120 mb. Thus, the mean-free-pass of a ^{3}He or t inside the target nucleus is small compared to size of a nucleus even for ^{12}C, $l = 1/n\sigma_{tot} \approx 0{,}5$ fm. It means that the target is black for the ^{3}He and t, and the charge-exchange takes place in a thin ridge on the surface of a nucleus. The distortion of an incoming and outgoing wave is mainly due to absorption and it can be taken into account in optical approximation. The wave function of an incoming ^{3}He is

$$\psi^{(+)}_{p_{He}}(\vec{r}) = \exp\left[i\vec{p}_{He}\cdot\vec{r} + \frac{2\pi i}{p_{He}} f(0) \int_{-\infty}^{z} \varrho(\vec{r}_\perp, z')\,dz'\right], \qquad (7)$$

and for outgoing tritium

$$\psi^{(-)}_{p_t}(\vec{r}) = \exp\left[i\vec{p}_t\cdot\vec{r} - \frac{2\pi i}{p_t} f^*(0) \int_{z}^{\infty} \varrho(\vec{r}_\perp, z')\,dz'\right] \qquad (8)$$

For heavy nucleus this approximation is reasonable due to small mean-free-pass of a projectile.

Now, one must use the wave functions (7),(8) in (5) instead of plane waves. Neglecting the change in q^2 due to admixture of small transversal components ($\sim 1/R$) one obtains

$$\frac{d^2\sigma}{d\Omega dE_t} = \frac{1}{2\pi}\frac{p_t}{p_{He}}\left(\frac{m_{He}}{2\pi m}\right)^2 g^2_N(q^2)\, F^2(q^2)\, q^2 \,\mathrm{Im}\, G^{dist}_\pi(q,q), \qquad (9)$$

where $G^{dist}_\pi(q,q)$ is

$$G^{dist}_\pi(q,q) = \int d^3r\, d^3r'\, G_\pi(\vec{r},\vec{r}')\, e^{-iq(r-r') - \frac{1}{2}\chi(r_\perp) - \frac{1}{2}\chi(r'_\perp)} \qquad (10)$$

The distortion factor is $e^{1/2\chi(r_\perp)}$ with

$$\chi(\vec{r}_\perp) = \sigma_{tot} \int_{-\infty}^{\infty} \varrho(\vec{p}_\perp, z)\,dz. \qquad (11)$$

The integration over z and z' in (10) can be performed explicitly.

$$G^{dist}_\pi(q,q) = \int d^2r\, d^2r'\, G_\pi(q_L, q_L; \vec{r}, \vec{r}')\, e^{-iq_\perp(r-r') - \frac{1}{2}\chi(r) - \frac{1}{2}\chi(r')} \qquad (12)$$

In the limit $q_L R \gg 1$, which is good approximation in the Δ region, the distorted function $G^{dist}_\pi(q,q)$ can be expressed in terms of pion propagator in infinite nuclear matter (see Appendix) The final expression for the cross-section is

$$\frac{d^2\sigma}{d\Omega dE_t} = \frac{p_t}{p_{He}}\left(\frac{m_{He}}{\pi\mu}\right)^2 f^2_N(q^2)\, F^2(q^2)\, q^2 \,\mathrm{Im}\, G_\pi(q,\omega;R)\cdot r\cdot b\cdot l/\hbar c, \qquad (13)$$

μ is the mass of a pion, f_N—πNN constant. The other notations see Appendix

Now, to calculate tritium spectrum we need to know pionic self-energy in the medium.

3. PIONIC SELF-ENERGY IN THE MEDIUM.

At low energy the pionic self-energy in nuclear matter has been studied in number of papers in connection with the pion-condensation [9-11]. In the Δ-region relativistic calculations are necessary. Below we shall follow the model developped in [12], with some modifications discussed later.

First, we would like to remind the outlines of the model. The change of the pionic self-energy in the medium comes from two sources. The main correction arises from the modification of nucleon propagator in the matter due to Pauli blocking.

$$\delta\hat{G}(p) = 2\pi i n_p \frac{\hat{p}+m}{2E_p}\,\delta(p_0 - E_p). \qquad (14)$$

It brings into play an amplitude of elastic pion-nucleon scattering with nucleons being on-shell and pions being off-shell. We are using the model developped in [13] for description both resonant and nonresonant partial waves of pion-nucleon amplitude near and below the Δ-resonance.

Another source of modifications is the short-range correlations in nuclear matter and nuclear absorption of virtual pions. The first one affects intermediate nucleon and isobar states by quenching their contributions. Here we are using the effective interaction in slightly different form than in [12]

$$W_{ph}=\frac{f_N^2}{\mu^2}g_N'(\vec{\tau}_1\cdot\vec{\tau}_2)(\vec{\sigma}_1\cdot\vec{\sigma}_2)+\frac{f_\Delta^2}{\mu^2}g_\Delta'(\vec{S}_1^+\cdot\vec{S}_2)(\vec{T}_1^+\cdot\vec{T}_2)+$$

$$+\frac{f_N\cdot f_\Delta}{\mu^2}g_{N\Delta}[(\vec{S}_1^+\cdot\vec{\sigma}_2)(\vec{T}_1^+\cdot\vec{\tau}_2)+\text{h.c.}] \tag{15}$$

letting the constants g_N' and g_Δ' be different.

Again, we would like to remind that the form (15) of the interaction is meaningful only when the exchange contributions are neglected.

The nuclear absorption is taken into account in a phenomenological way by using the imaginary part of the pion optical potential [14,15], which describes the p-wave absorption at small q. Following this, we add the term

$$4\pi i C_0 n^2 q^2, \tag{16}$$

with $C_0=0.08\left(\frac{\hbar}{\mu c}\right)^6$ to the pion self-energy.

The final expression for the pion self-energy is

$$\Pi(\omega,\vec{q})=\Pi_N(\omega,\vec{q})+\Pi_\Delta(\omega,\vec{q})+\Pi_\sigma(\omega,\vec{q})+4\pi i C_0 n^2\vec{q}^2, \tag{17}$$

where

$$\Pi_N(\omega,\vec{q}) = q^2\tilde{\chi}_N(\omega,\vec{q}),$$

$$\Pi_\Delta(\omega,\vec{q}) = \frac{4}{9}\left(\frac{f_\Delta(q^2)}{\mu}\right)^2\frac{m}{m_\Delta^2}n\vec{q}^2+\frac{8}{9}\left(\frac{f_\Delta(q^2)}{\mu}\right)^2\frac{1}{m_\Delta}\left(1+\frac{m}{m_\Delta}\right)n\cdot q^2+$$

$$+\frac{4}{9}\frac{f_\Delta^2(q^2)}{\mu^2}\frac{(m+m_\Delta)^2-q^2}{m_\Delta^2}\vec{q}^2\tilde{\chi}_\Delta(\omega,\vec{q}). \tag{18}$$

and Π_σ arises from the sigma-term of pion-nucleon amplitude,

$$\Pi_\sigma(\omega,\vec{q})=\frac{1}{f_\pi^2}\left(1-\frac{2q^2}{\mu^2}\right)\sigma(0)\cdot n, \tag{19}$$

where f_π is pion decay constant, $f_\pi=133$ MeV and $\sigma(0)$ is the value of σ-commutator for forward scattering; $\sigma(0)=66$ MeV, $\tilde{\chi}_N$ and $\tilde{\chi}_\Delta$ are nucleon and Δ spin susceptibilities corrected for short-range correlations.

$$\tilde{\chi}_N=\chi_N\frac{1+\chi_\Delta\cdot(g_{n\Delta}-g_\Delta)}{(1-g_N\cdot\chi_N)\cdot(1-g_\Delta\chi_\Delta)-g_{N\Delta}^2\cdot\chi_N\cdot\chi_\Delta}, \tag{20}$$

$$\tilde{\chi}_\Delta=\chi_\Delta\frac{1+\chi_N\cdot(g_{n\Delta}-g_N)}{(1-g_N\cdot\chi_N)\cdot(1-g_\Delta\cdot\chi_\Delta)-g_{N\Delta}^2\cdot\chi_N\cdot\chi_\Delta}. \tag{21}$$

The expressions for χ_N and χ_Δ are [12]

$$\chi_N(\omega,\vec{q})=4\frac{f_N^2(q^2)}{\mu^2}\int\frac{d^3p}{(2\pi)^3}n_{\vec{p}}\left[\frac{1}{\omega-\frac{\vec{p}\cdot\vec{q}}{m}+\frac{q^2}{2m}}+\right.$$

$$\left.+\frac{1}{-\omega+\frac{\vec{p}\cdot\vec{q}}{m}+\frac{q^2}{2m}}+2\pi i n_{\vec{p}+\vec{q}}\,\delta(\omega+\varepsilon_{\vec{p}}-\varepsilon_{\vec{p}+\vec{q}})\right]; \tag{22}$$

with $q^2=\omega^2-\vec{q}^2$; and

$$\chi_\Delta(\omega,\vec{q}) = \frac{4}{9}\frac{f_\Delta^2(q^2)}{\mu^2}\int\frac{d^3p}{(2\pi)^3}n_{\vec{p}}\left[\frac{1}{\omega-\frac{\vec{p}\cdot\vec{q}}{m}+\frac{q^2}{2m}-\frac{(\tilde{m}_\Delta^2-m^2)}{2m}}+q\leftrightarrow -q\right], \tag{23}$$

where $\tilde{m}_\Delta^2=m_\Delta^2-im_\Delta\Gamma_\Delta(\omega,\vec{q})$.

The properties of the Δ itself, including the width $\Gamma_\Delta(\omega,\vec{q})$, remain unchanged in our approach. For the Γ_Δ we have

$$\Gamma_\Delta(\omega,\vec{q})=\Gamma_0\left(\frac{k}{k_0}\right)^3\frac{k_0^2+\varkappa^2}{k^2+\varkappa^2}, \tag{24}$$

where k is the relative momentum of a pion and a nucleon in c.m.

$$k^2=\frac{(M^2+m^2-\mu^2)^2}{4M^2}-m^2, \tag{25}$$

with M being a current mass $M^2=m^2+2m\omega+q^2$, and $k_0=k(M=m_\Delta)$. $\Gamma_0=120$ MeV and $\varkappa=200$ MeV/c. The parameter of centrifugal barrier $\varkappa$ is necessary to reproduce the shape of a resonance in $\pi N\to\pi N$ reaction.

Strictly speaking, the expressions (24) for Γ_Δ and (17) for Π_Δ are not consistent with each other. Im Π_Δ is proportional to elastic pion-nucleon cross-section, i. e. Im $\Pi_\Delta\sim\Gamma_\Delta^2$. To do Π_Δ and Γ_Δ consistent one must include into the vertexes in Π_Δ the factor $\sqrt{\frac{k_0^2+\varkappa^2}{k^2+\varkappa^2}}$ where k^2 depends on a virtual pion mass q^2 instead of μ^2.

A pion spectrum with the self-energy (17) is shown in Fig. 2 for two sets of constants g'. The spectrum is mostly sensitive to the constant g'_Δ. An imaginary part of the pion propogator calculated for $g'_N = g'_{N\Delta} = 0$, $g'_\Delta = 0.5$ is shown in Fig. 3. At small $\vec{q}$ the pionic peak dominates completely in Im $G_\pi(\omega, \vec{q})$, but at higher q, entering into Δ-hole continuum, it dissappears due to strong mixing with Δ-hole excitations. Another feature that is seen from Fig. 3 is the second maximum at high omega corresponding to the «true» Δ-hole collective mode.

4. RESULTS AND DISCUSSIONS

Using the expression (13) for the cross-section and the pion self-energy (17) we calculated tritium spectrum for forward charge-exchange at different energies of incoming ^{3}He.

The tritium spectrum from the (^{3}He, t) reaction on carbon for $\theta = 0^0$ at the kinetic energy of ^{3}He [4] $E = 2.3$ GeV is shown in Fig. 4. The cross-sections $\frac{1}{p_t}\frac{d^2\sigma}{dE_t d\Omega}$ measured in [3] for $\theta = 0^0$ at the momentum of ^{3}He $p = 4399.7$ MeV/c is shown in Fig. 5. In both cases the position and height of the peak is reproduced satisfactorily in calculations. At higher energies, however, the agreement is much more poor. The cross-sections $\frac{1}{p_t}\frac{d^2\sigma}{dE_t d\Omega}$ at the momenta of ^{3}He $p = 6807.1$ MeV/c and $p = 10790.3$ MeV/c are shown in Fig. 6 and Fig. 7. At the momentum of ^{3}He $p = 6807.1$ MeV/c the position of observed peak is higher comared with the predictions and its width is larger for factor ~ 2. At the momentum $p = 10790.3$ MeV/c the calculated cross-section comprise only a small portion of the observed one.

There could be several reasons for such a strong disagreement. The obvious one is that the approximations used in deriving (13) for the cross-section can not be justified at high energy. In particular, the condition $qR \gg 1$ is apparently violated at the momentum $p_{He} = 10790$ MeV/c for carbon target. The finite size effects shift the peak position to higher energies [16]. It would be very interesting to repeat the measurements at high energy for heavy nuclei to separate these effects.

From the other hand, the definite increase of the width indicates the existence of other mechanisms of energy absorption for charge-exchange on a nucleus which absent on a single nucleon. Partially, these mechanisms are taken into account by an imaginary

optical potential in (17). However, this is not sufficient to reproduce the observed width.

In general, one can conclude that for comparatively small kinetic energy of ^{3}He, when $qR \gg 1$, the proposed theory reproduces satisfactory the effects observed in charge-exchange reaction at high excitation energy of a nucleus. At higher energy more accurate account of finite size effects is necessary before discussion of other reaction mechanisms.

In conclusion, I would like to acknowlege V.G. Ableev, C. Gaarde, E.A. Strokovsky, L.N. Strunov, O.P. Sushkov, V.I. Inozemtsev and S.M. Eliseev for numerous discussions during the work.

APPENDIX.

Let us rewrite the Eq. (10) introducing new coordinates $\vec{\varrho}=\vec{r}-\vec{r}', \ \vec{R}=\frac{\vec{r}+\vec{r}'}{2}$,

$$G_\pi^{\text{dist}}(q,q)=\int d^3\varrho\, d^3R\, G_\pi(\vec{\varrho},\vec{R})\, e^{-iq\varrho-\frac{1}{2}\chi\left(\vec{R}_\perp+\frac{\varrho_\perp}{2}\right)-\frac{1}{2}\chi\left(\vec{R}_\perp-\frac{\varrho_\perp}{2}\right)}. \tag{A1}$$

integration over logitudinal coordinate ϱ_L (along $\vec{p}_{\text{He}}$) can be performed explicitly

$$G_\pi^{\text{dist}}(q,q)=\int d^2\varrho\, d^3R\, G_\pi(q_L;\vec{\varrho},\vec{R})\, e^{-iq_\perp\varrho-\frac{1}{2}\chi\left(\vec{R}_\perp+\frac{\varrho}{2}\right)-\frac{1}{2}\chi\left(\vec{R}_\perp-\frac{\varrho}{2}\right)}. \tag{A2}$$

Let us now introduce a Fourie-transform of G_π over the transversal coordinate $\vec{\varrho}$. Then we obtain

$$G_\pi^{\text{dist}}(q,q)=\int d^3R\, d^2\varrho\,\frac{d^2\varkappa}{4\pi^2}\, G_\pi(q_L,\vec{\varkappa};\vec{R})\, e^{-i(q_\perp-\varkappa)\varrho-\frac{1}{2}\chi\left(\vec{R}_\perp+\frac{\varrho}{2}\right)-\frac{1}{2}\chi\left(\vec{R}_\perp-\frac{\varrho}{2}\right)}. \tag{A3}$$

The integration over $\vec{\varrho}$ in (A3) goes outside the region $|\vec{\varrho}|\approx R_0$ where R_0 is a nuclear radius. At distances smaller than that $\chi(\vec{r})$ is large and contribution into the integral is suppressed by the absorption factor $e^{-\chi(\vec{r})}$. In this case the integral over $\vec{\varrho}$ as a function of a $\vec{\varkappa}$ has a sharp maximum at $\vec{\varkappa}=\vec{q}_\perp$ having the width $\sim 1/R_0$. At $q_L R_0 \gg 1$ the function $G_\pi(q_L,\vec{\varkappa};\vec{R})$ as a function of $\vec{\varkappa}$ is smooth inside the maximum and can be taken at $\vec{\varkappa}=\vec{q}_\perp$ and moved outside the integral. The remaining integration over $\vec{\varkappa}$ and $\vec{\varrho}$ is trivial and we obtain:

$$G_\pi^{\text{dist}}(q,q)=\int d^3R\, G_\pi(\vec{q};\vec{R})\, e^{-\chi(\vec{R}_\perp)}. \tag{A4}$$

The main contribution into (A4) for Im G_π^{dist} comes from the region near nuclear surface where $\chi(R)$ is small. Near the surface, for $q_0 R_0 \gg 1$ the Im $G_\pi(q,R)$ follows nuclear density profile, therefore

$$\text{Im}\, G_\pi^{\text{dist}}(q,q)\approx \text{Im}\, G_\pi(\vec{q};R_0)\,\frac{1}{n\sigma_{\text{tot}}}\int d^2R_\perp\,\chi(R_\perp)\, e^{-\chi(\vec{R}_\perp)}. \tag{A5}$$

Near the surface the density falls down linearly with increasing R. Thus, we obtain

$$\int d^2R_\perp\,\chi(R_\perp)\, e^{-\chi(R_\perp)}\approx 2\pi R_0 b. \tag{A6}$$

where b is the width of a region where the density decreases linearly down to zero. For Fermi-function of density radial dependence $n(r)\sim 1/\left(\exp\left(\frac{r-R_0}{\alpha}\right)+1\right)$, $b\approx\pi\alpha$. Thus, for $qR_0\gg 1$ we have

$$\text{Im}\, G_\pi^{\text{dist}}(q,q)\approx \text{Im}\, G_\pi(q;R)\, 2\pi l b R, \tag{A7}$$

where $l=1/\pi\sigma_t$.

REFERANCES

1. *C. Goodman.* Nucl. Phys. A374 (1982) 241.
2. *C. Gaarde.* Nucl. Phys. A396 (1983) 127.
3. *V.G. Ableev et al.*, Pisma v JETF, 40 (1984) 35.
4. *C. Gaarde.* The talk on Niels Bohr Centennial Conference, May, 1985.
5. *T.E.O. Ericson, J. Hüfner.* Phys. Lett., 33B (1970) 601.
6. *C.B. Dover and R.H. Lemmer.* Phys. Rev., C7 (1973) 2312.
7. *H. Esbensen, H. Toki and G.F. Bertsch.* Phys. Rev., C31 (1985) 1816.
8. *H. Esbensen, and T.-S.H. Lee.* Phys. Rev., C32 (1986) 1966.
9. *A.B. Migdal.* Rev. Mod. Phys., 50 (1978) 107.
10. *W.M. Alberico, M. Ericson and A. Molinari.* Phys.Lett., 92B (1980) 15.
11. *I.N. Borzov, E.E. Saperstein, S.V. Tolokonnikov and S.A. Fajans.* Particles and Nuclei, 12 (1981) 848.
12. *.F. Dmitriev and Toru Suzuki.* Nucl. Phys., A438 (1985) 697.
13. *M.G. Olsson and E.T. Osypovsky.* Nucl. Phys., B101 (1975) 136.
14. *E. Oset, W. Weise and R. Brockman. Phys. Lett., 82B (1979) 344.*
15. *W.M. Alberico, M. Ericson and A. Molinari. Phys. Lett., 136B (1984) 307.*
16. *G. Chanfray and M. Ericson*, Phys. Lett. 141B (1984) 163.

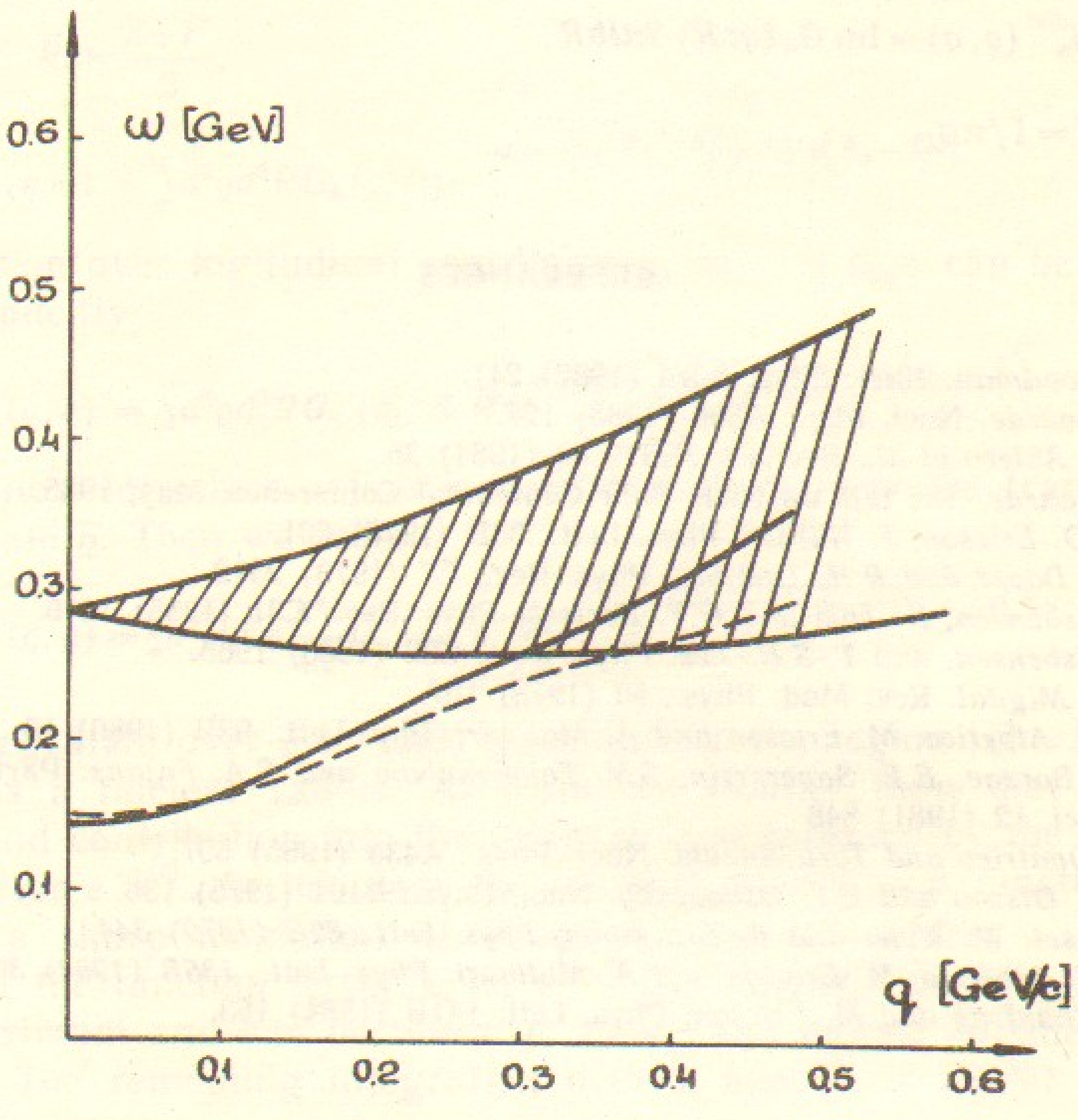

Fig. 2. Pionic mode in nuclear matter. Dashed line is the spectrum without short-range repulsion. Full line corresponds to $g'_N = g'_{N\Delta} = 0$, $g'_\Delta = 0.5$.

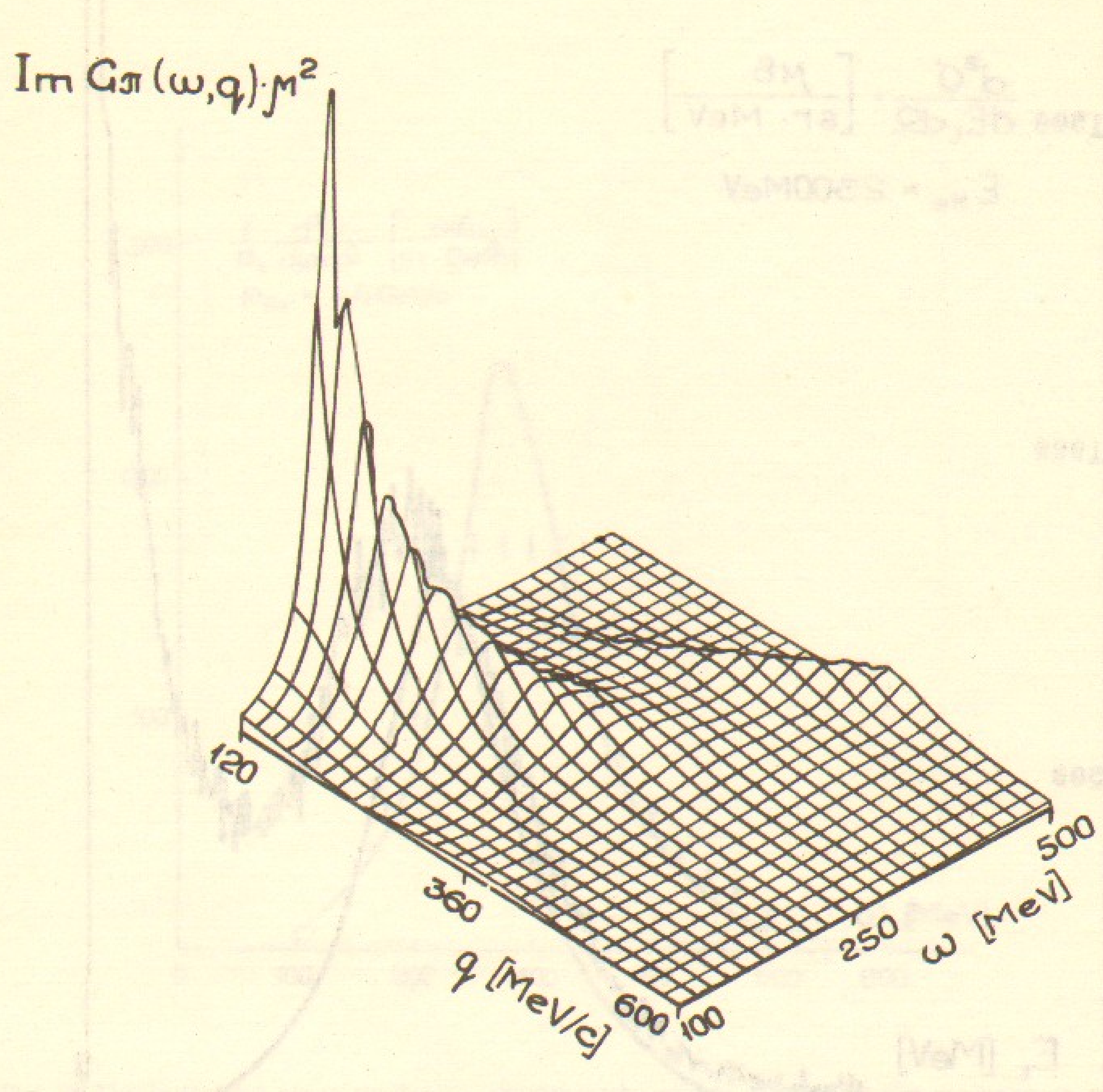

Fig. 3. The shape of imaginary part of the pion propogator in units of inverse pion mass squared for the constants of short-range repulsion $g'_N = g'_{N\Delta} = 0$, $g'_\Delta = 0.5$.

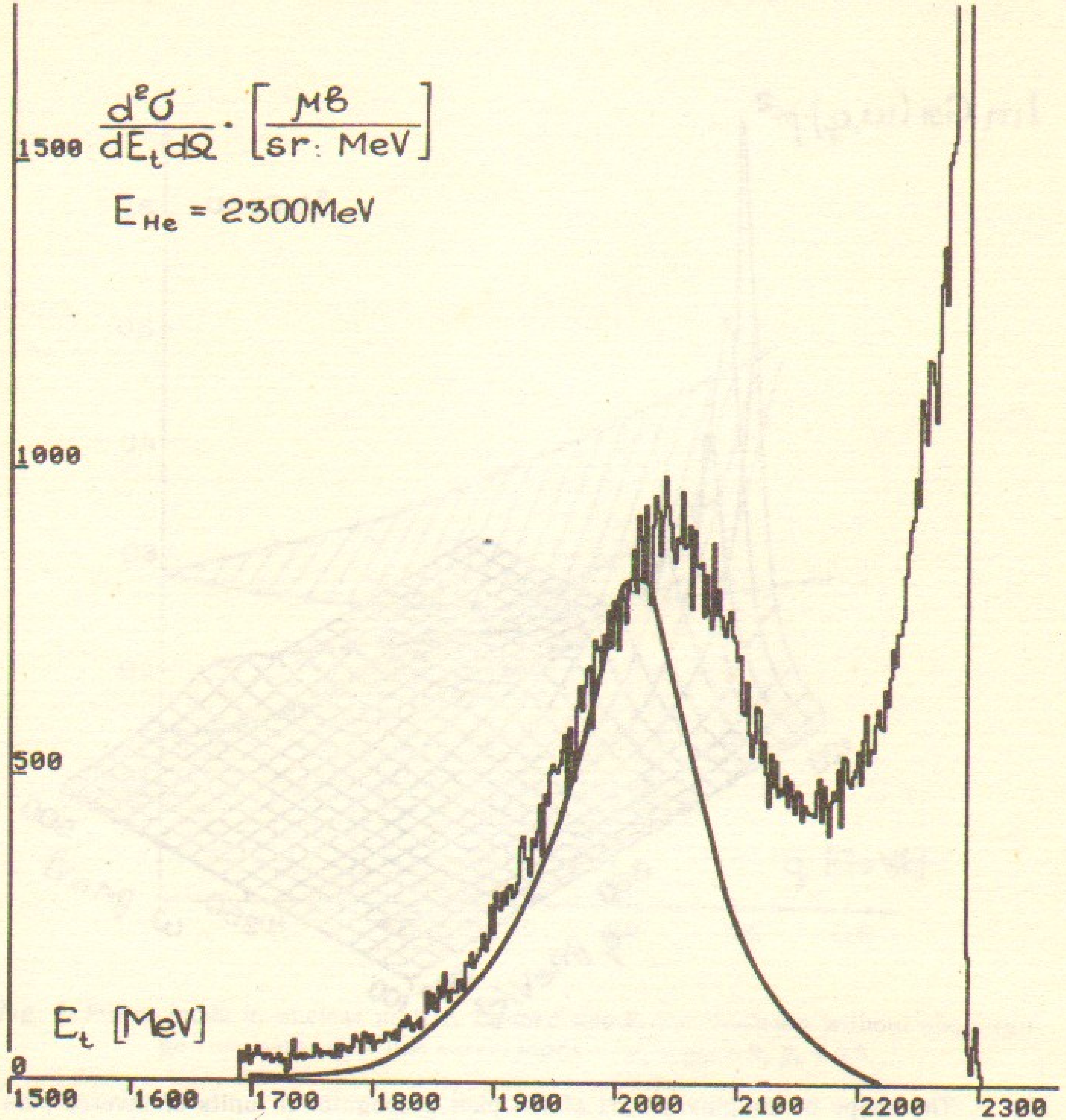

Fig. 4. Tritium spectrum from charge-exchange reaction (^{3}He, t) on carbon [4] at the kinetic energy of ^{3}He $E=2300$ MeV. The theoretical curve corresponds to . $g'_N = g'_{N\Delta} = 0$, $g'_\Delta = 0.5$.

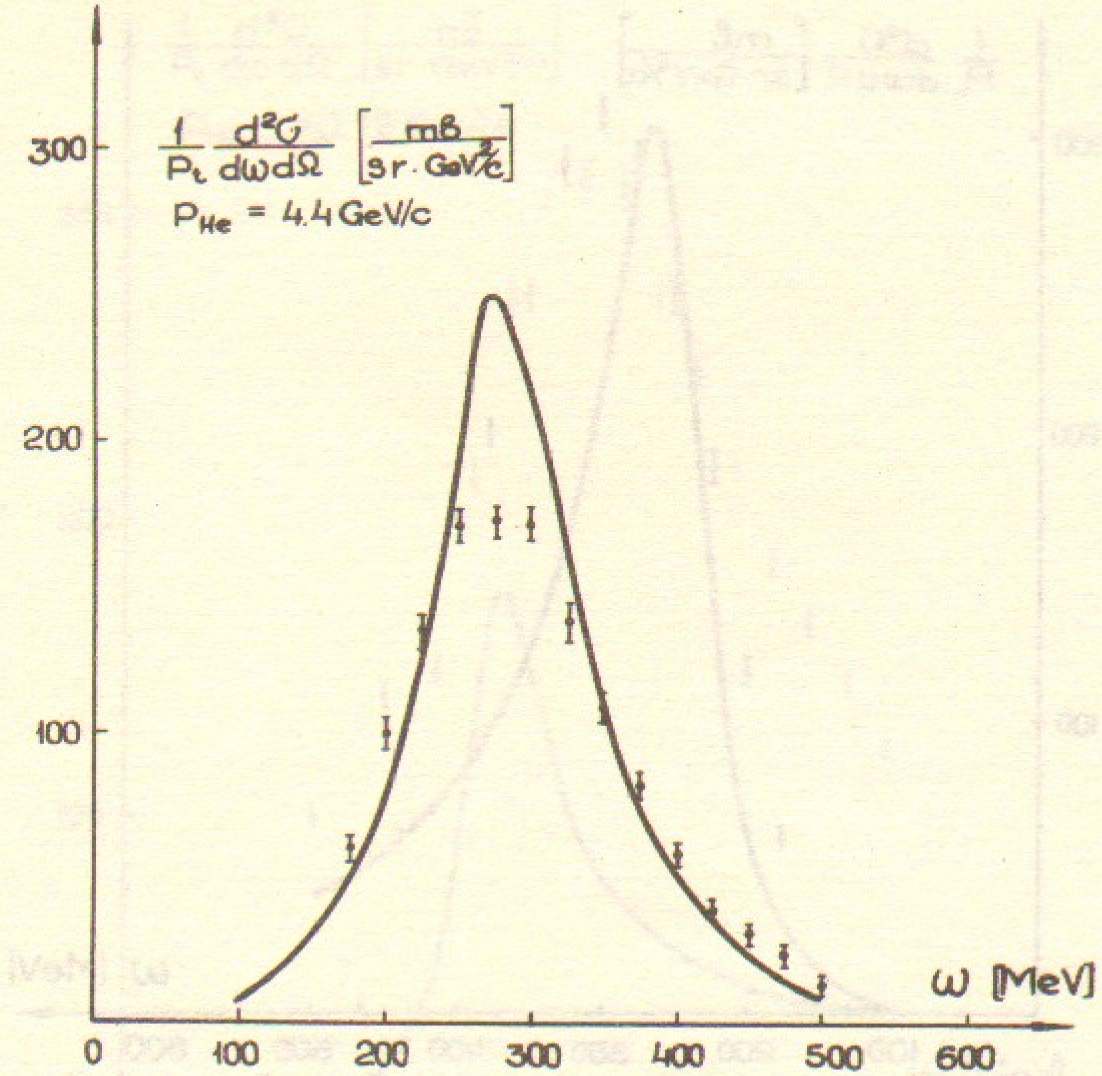

Fig. 5. The invariant cross-section of charge-exchange on carbon at the momentum of ^{3}He $p=4.4$ GeV/c. Theoretical curve is taken with the same constants as in Fig. 4.

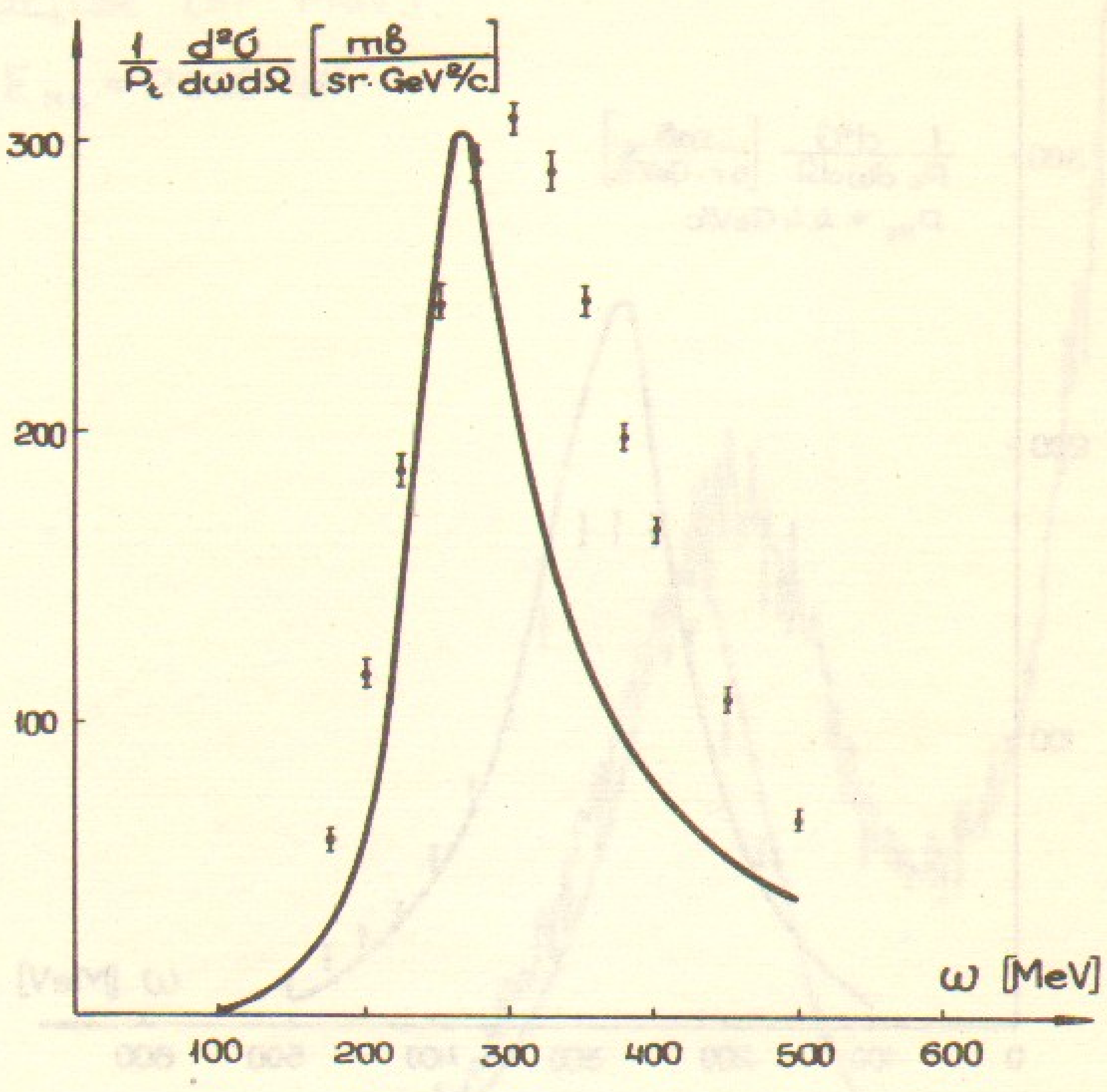

Fig. 6. The invariant cross-section of charge-exchange on carbon at the momentum of ^{3}He $p=6.79$ GeV/c. Theoretical curve is taken with the same constants.

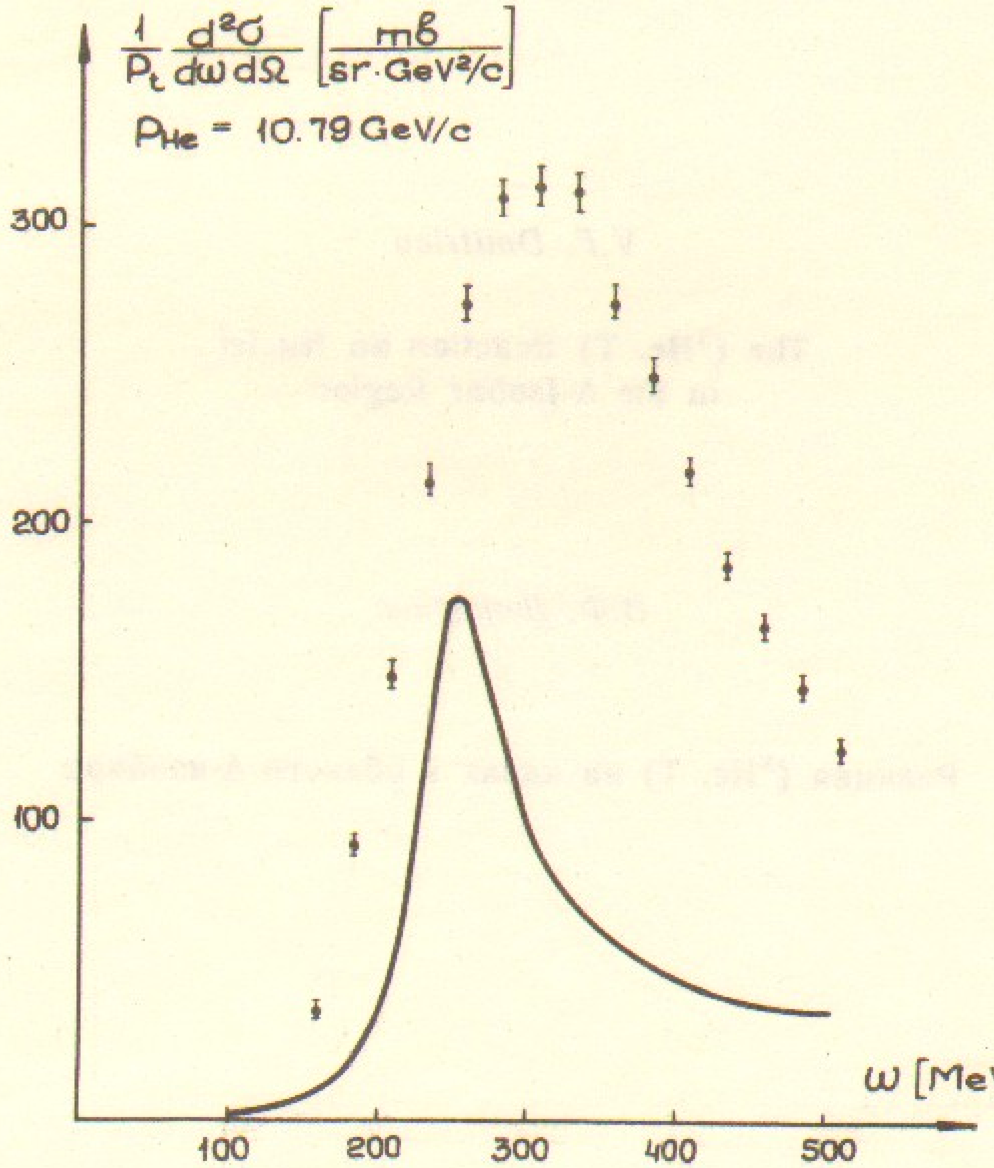

Fig. 7. The invariant cross-section of charge-exchange on carbon at the momentum of ^{3}He $p=10.79$ GeV/c. Theoretical curve is taken with the same constants.

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The (^{3}He, T) Reaction on Nuclei in the Δ-Isobar Region

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Реакция (^{3}He, T) на ядрах в области Δ-изобары

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