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ON INELASTIC SCREENING IN K_S REGENERATION ON NUCLEI

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A B S T R A C T

We discuss the inelastic screening corrections to the kaon-nucleon total cross-sections and K_S regeneration on nuclei in the framework of the parton model of diffraction scattering.

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Recent high precision data ^{1),2)} on K_L nucleus interactions has once more raised the issue of inelastic screening in particle-nucleus interactions. First introduced by Gribov ³⁾, inelastic screening is due to the inelastic, $h^* \neq h$, intermediate states in the diagrams of Fig. 1. Inelastic screening is definitely vital to reproduce the total cross-sections ⁴⁾. Recent data on $\sigma_{\text{tot}}(K_L A)$ perfectly fit the expected pattern : measured cross-sections are lower than those given by the simple optical model and allowance for inelastic screening brings theory and experiment to comfortable agreement ⁵⁾.

What came out as a surprise was the behaviour of the $K_L A \rightarrow K_S A$ regeneration amplitude f_{LS}^A . The effective intercept of the ω trajectory, " $\alpha_\omega(0)$ " = 0.39 ± 0.01 , differs drastically from $\alpha_\omega(0) = 0.44 \pm 0.01$, as was found in kaon-nucleon interactions. Moreover, including the inelastic screening corrections was found to worsen the description of the magnitude of the regeneration amplitude ⁵⁾.

Bertocchi and Treleani were the first to demonstrate the importance of the inelastic intermediate states, produced via $C = -1$ exchanges in the t channel ⁶⁾. Their approach was further developed by Diu and Ferraz De Camargo ⁷⁾, who, at the price of a number of assumptions and of few adjustable parameters, have produced rather good fit to both $\sigma_{\text{tot}}(K_L A)$ and f_{LS}^A .

Still, we find the situation unsatisfactory. Starting from the paper of Karmanov and Kondratyuk ⁸⁾, an assumption

$$F_{K^*N} = F_{KN} \quad (1)$$

is widely, and uncritically, used for estimates of screening corrections. However, this assumption is grossly inconsistent. It has long been well known that, if equality (1) holds, there would be no inelastic diffraction scattering altogether ⁹⁾. In the formalism of Ref. 8), applied in Refs. 6) and 7) to K_L nucleus interactions, all connection between the underlying mechanisms of inelastic diffraction scattering and of inelastic screening is lost.

In this paper we propose a novel description of the screening corrections in regeneration amplitude. Our goal is not to produce yet another good fit to the data of Refs. 1) and 2). Rather we demonstrate explicitly how the screening corrections come out from the explicit model, incorporating inelastic diffraction scattering ¹⁰⁾, and how inelastic screening allows us to determine some of the parameters which enter into the modern parton model interpretation of Reggeon field theory ¹¹⁾.

We begin by reviewing the modern multichannel theory of diffraction scattering ^{9),10),12)}. Wave functions $|\psi\rangle$ of hadrons are decomposed in terms of eigenstates $|m\rangle$ of the scattering matrix :

$$|\psi\rangle = \sum_m C_m |m\rangle \quad (2)$$

The usual orthogonality, $\langle i|k\rangle = \delta_{ik}$, and completeness, $\sum_m |C_m|^2 = 1$, conditions are assumed. The final state of scattering,

$$\hat{F}|\psi\rangle = \sum_m C_m f_m |m\rangle \quad (3)$$

does not coincide with the initial state, provided that not all the scattering amplitudes f_m are identical. In the framework of the modern parton model the scattering eigenstates are identified with states with a fixed number of soft partons ¹²⁾. An important observation, due to Grassberger ¹¹⁾, is that there should be eigenstates containing no soft partons at all. Such an eigenstate, to be called hereafter the passive state, has a vanishing interaction cross-section. The weight $P = |C_0|^2$ of the passive component remains unspecified. However, as proved in Ref. 11), P is a rising function of energy, approaching some asymptotic value.

An integral inelastic diffraction cross-section is given by equation ¹²⁾ :

$$\sigma_{diff} = \int d^2\vec{b} \sum_m |C_m|^2 (f_m^2 - \langle f \rangle^2) \quad (4)$$

where $\langle f \rangle = \sum_m |C_m|^2 f_m$. There are different sources of fluctuations of f_m around $\langle f \rangle$ ¹²⁾. First, there are fluctuations in the number n of soft partons. Second, for a fixed number of soft partons there are impact parameter $\vec{b}_i$ fluctuations. For interactions with the nuclei the latter

fluctuations can safely be neglected, and it suffices to consider only the fluctuations in the number of soft partons.

An obvious generalization of Eq. (3) to the particle-nucleus interactions is ^{10),13)}

$$\sigma_A = 2 \int d^2\vec{b} \sum_m |C_m|^2 (1 - \exp[-\sigma_m T(\vec{b})/2]) \quad (5)$$

whereas the usual optical formula reads ^{*)} :

$$\sigma_A = 2 \int d^2\vec{b} (1 - \exp[-\langle\sigma\rangle T(\vec{b})/2]) \quad (6)$$

Here

$$\langle\sigma\rangle = \sum_m |C_m|^2 \sigma_m ; \quad T(\vec{b}) = \int dz \rho(z, \vec{b})$$

is the usual profile function. It should be emphasized that Eq. (5) corresponds to the exact summation of all the multiple scattering diagrams of Fig. 1 ¹⁰⁾. Due to the inequality

$$\langle \exp[-\frac{\sigma}{2} T(\vec{b})] \rangle \geq \exp[-\frac{\langle\sigma\rangle}{2} T(\vec{b})] \quad (7)$$

an inelastic screening reduces the total cross-sections. For the difference of K^0A and $\bar{K}^0A$ cross-sections we obtain an equation

$$\sigma_A - \bar{\sigma}_A \cong \int d^2\vec{b} T(\vec{b}) \sum_m |C_m|^2 (\sigma_m - \bar{\sigma}_m) \exp[-\frac{\sigma_m + \bar{\sigma}_m}{4} T(\vec{b})] \quad (8)$$

which should be compared with its usual counterpart :

$$\sigma_A - \bar{\sigma}_A \cong (\sigma_N - \bar{\sigma}_N) \int d^2\vec{b} T(\vec{b}) \exp[-\frac{\sigma_N + \bar{\sigma}_N}{4} T(\vec{b})] \quad (9)$$

^{*)} We refer to the optical model only for the sake of simplicity. In all the calculations below, instead of $\exp[-\sigma T(\vec{b})/2]$ we use $(1 - \sigma T(\vec{b})/2A)^A$ ¹⁴⁾.

Let us now formulate a model ¹⁰⁾ which, though simple, incorporates all the essential features of any sophisticated theory. We consider hadrons as made of constituent quarks. To each constituent quark we assign, with probability P_q , a passive state, and with probability $A_q = 1 - P_q$, an active state. It is the presence of the passive component, which gives the principal contribution to inelastic screening, so that we neglect the fluctuations within the active component. In such a model $\sigma_{qN} = A_q \sigma$, and

$$\sigma_{KN} = (A_u + A_s) \sigma = A_u (1 + \lambda) \sigma \quad (10)$$

where $\lambda \simeq 0.5$, and σ stands for the interaction cross-section for the active component. The resulting equations for the nuclear cross-section are of the form :

$$\begin{aligned} \sigma_{\text{tot}}(K_L A) = 2 \int d^2 \vec{b} \left\{ \lambda A_u^2 (1 - \exp[-\sigma T(\vec{b})]) \right. \\ \left. + [\lambda A_u (1 - A_u) + A_u (1 - \lambda A_u)] (1 - \exp[-\sigma T(\vec{b})/2]) \right\} \end{aligned} \quad (11)$$

$$\begin{aligned} \text{Im } f_{LS}^A = \text{Im } f_{LS}^N \int d^2 \vec{b} T(\vec{b}) \left\{ \lambda A_u \exp[-\sigma T(\vec{b})] \right. \\ \left. + (1 - \lambda A_u) \exp[-\sigma T(\vec{b})/2] \right\} \end{aligned} \quad (12)$$

In the last equation we have taken into account that in the quark model the s quark decouples from the ω trajectory. Actually, $\text{Im } f_{LS}^N$ in Eq. (12) should be understood as $\text{Im } f_\omega + (1 - 2Z/A) \text{Im } f_\rho$, where f_ω and f_ρ are the ω and ρ exchange contributions.

According to Grassberger's general theorem $dA_u/dy = -\gamma < 0$ ¹¹⁾ [hereafter $y = \ln(s/m_p^2)$]. The magnitude of γ is controlled by the gap $\Delta = \alpha_p(0) - 1$. For instance, in the model of Ref. 15),

$$\gamma = \Delta (A_u(y))^2 e^{-\Delta y} / A_u(\infty) \sim \Delta \quad (13)$$

Equation (13) cannot be applied literally to our over-simplified model. Rather we take γ and $A_0 = A_u(y=5)$ as two free parameters of our model and use an expansion $A_u(y) = A_0 - \gamma(y-5)$. For the KN total

cross-section and f_{LS}^N we use the same parametrization as was used in Ref. 6). Equation (12) gives the imaginary part of the regeneration amplitude. The real part of f_{LS}^A and the regeneration phase are calculated using the derivative analyticity relations¹⁶⁾.

$\sigma_{tot}(K_L N)$ is known to be a rising function of energy. Since $A_u(y)$ decreases at high energies, σ rises even faster than $\sigma_{tot}(K_L N)$. Therefore, the $1 - \exp[-\sigma T(b)]$ and $1 - \exp[-\sigma T(b)/2]$ factors in the integrand of Eq. (11) are rising functions of energy. As a result, for the light nuclei $\sigma_{tot}(K_L A)$ also rises. However, for the black nuclei, $\exp[-\sigma T(b)] \simeq 0$, so that $\sigma_{tot}(K_L A) \simeq 2\pi R_A^2 A_u(y)(1+\lambda)$ should, instead, be decreasing. There is certain evidence for such behaviour of $\sigma_{tot}(K_L A)$, but the accuracy of the data is not sufficient to draw a definite conclusion. In Fig. 2, we present $\sigma_{tot}(K_L A)$, given by Eq. (11) with $A_0 = 0.8$, versus the experimental data of Ref. 1). These data allow us to determine A_0 , but not the value of γ .

Having A_0 specified, we turn now to the regeneration data. Let us make a finite y expansion $\sigma_{tot}(K_L N) = \sigma_0 + \sigma_1 y$. Then we get

$$\exp[-\sigma T(\vec{b})] \simeq \exp[-\sigma_0 T(\vec{b})/A_0(1+\lambda)] S^{-\delta} \quad (14)$$

where

$$\delta = (\sigma_1 + \sigma_0 \gamma/A_0) T(\vec{b})/A_0(1+\lambda) \quad (15)$$

Thus, we arrive at the atomic number dependent renormalization of the intercept of the ω trajectory :

$$"\alpha_\omega(0)" = \alpha_\omega(0) - \langle \delta \rangle \quad (16)$$

The sign of the renormalization is in perfect agreement with that observed experimentally. We emphasize that such a renormalization is quite natural in the supercritical Reggeon field theory with rising cross-sections¹⁷⁾, in which $\Delta = \alpha_P(0) - 1 > 0$. Equation (12) corresponds to an explicit summation of all the eikonal type ω multi-Pomeron exchanges. Intercepts of secondary trajectories are not renormalized only in the case of $\alpha_P(0) \equiv 1$. However,

$\Delta > 0$, the intercepts of secondary trajectories are no longer universal and depend on the target, as we have demonstrated explicitly. In Fig. 3, we present f_{LS}^A , calculated using Eq. (12) with $\gamma = 0.08$, versus the experimental data of Ref. 2). Since $\Delta \simeq 0.07$ ¹⁷⁾, this value of γ agrees well with that expected from Eq. (13). In Fig. 4, we present our predictions for the energy and atomic number dependence of the regeneration phase φ_{LS}^A . According to Refs. 18) and 19) in the momentum range $(30 \div 150)$ GeV/c $\varphi_{LS}^{pb} = -(122 \pm 1.8)^\circ$ and $\varphi_{LS}^C = -(125 \pm 1.6)^\circ$, in good agreement with our predictions. The precision of the available data ¹⁹⁾ on the regeneration phase for the carbon nucleus is not sufficient to test the predicted energy dependence of φ_{LS}^A . However, this variation of φ_{LS} seems to be of detectable magnitude for the lead nucleus.

In the limit of $\sigma T(b) \gg 1$ the integrands in Eqs. (8) and (9) are drastically different. Therefore, one might expect large deviations of the differential cross-section $d\sigma_{LS}^A/dt$ from that given by a simple optical model. Unfortunately, this is not the case. First, for a real nucleus $\sigma T(b) \sim 1$, so that differences between the profile functions in Eq. (8) and Eq. (9) are, to a great extent, washed out ^{*}). Second, in the region of diffraction minima coherent and incoherent regeneration cross-sections are of comparable magnitude ¹⁸⁾, so that very high precision measurements are needed to distinguish between the $d\sigma_{LS}^A/dt$, given by Eqs. (8) and (9). Third, for an actual calculation of the differential cross-section in the dip-bump region the impact parameter fluctuations are implicitly to be taken into account.

A few concluding remarks are in order. Our fits to the total cross-section and regeneration amplitude are not much better than, say, those of Ref. 7). However, our treatment of screening corrections has the unquestionable advantage of being based on the explicit model of diffraction scattering. Our conclusion is that inelastic screening corrections are definitely needed for a quantitative description of $\sigma_{tot}(K_L A)$ and f_{LS}^A . Second, the regeneration data give firm evidence for the specific energy dependence of the diffraction cross-section due to the energy dependence of

^{*}) For the lead nucleus a position of the first diffraction dip shifts to smaller values of $|t|$ by $\Delta t \simeq 2 \cdot 10^{-4} (\text{GeV})^2$, and the height of the first diffraction bump is 20% lower compared to those given by the simple eikonal model.

the fraction of the passive component. The local derivative σ_1 of the kaon-nucleon cross-section equals $\sigma_1 \simeq 1.8$ mb. Since $\sigma_0 \simeq 20$ mb, we have $\sigma_0 \gamma \simeq 2$ mb, so that the non-vanishing value of γ is indeed important quantitatively. Actually, $\gamma = 0.04$ is already grossly inconsistent with the data. Third, the value of γ , required by the experimental data, is in comfortable agreement with the value of γ , which is expected on quite general grounds, thus giving confidence on the consistency of the approach as a whole.

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FIGURE CAPTIONS

- Fig. 1 Multiple scattering diagrams for particle-nucleus scattering with inelastic intermediate states h^* .
- Fig. 2 Predicted kaon-nucleus total cross-sections versus the data of Ref. 1).
- Fig. 3 Predicted K_S regeneration amplitudes versus the data of Ref. 2).
- Fig. 4 Predicted momentum and atomic number dependences of K_S regeneration phase φ_{LS} .

FIG . 1

FIG. 2

FIG. 3

FIG. 4