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$\bar{K}^0 p$ and $\bar{p} p$ Elastic Scattering and
the Shadow of Generalized Annihilation Processes

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A B S T R A C T

The differences between the high energy total cross sections and elastic scatterings of $\bar{K}^-p$ and $\bar{p}p$ and their counterparts under crossing, K^+p or pp , are analysed in terms of the inelastic channels open for the former and closed for the latter systems. For $\bar{p}p$ these channels are mainly baryon annihilation. For $\bar{K}^-p$ we introduce the concept of hypercharge annihilation $\bar{K}^-p \rightarrow$ hypercharge zero particles (mainly one $Y = 0$ hyperon and $Y = 0$ mesons) and a vaguer notion of charge annihilation, and we propose that these generalized annihilation channels are the main ones open to $\bar{K}^-p$ and closed to K^+p . The contribution of the generalized annihilation processes to the overlap function of the overall collision is derived from the σ_T and $d\sigma_{el}/dt$ data, using various simple model assumptions. The results, which are insensitive to the latter, predict the total annihilation cross sections and their high energy dependence, as well as the range of the annihilation processes. The annihilation cross sections compare well with the few measured values. The annihilation ranges are predicted to be equal for $\bar{K}^-p$ and $\bar{p}p$, with a value close to the range of elastic scattering of $\bar{p}p$ and pp . The elastic range of Kp is smaller by about a factor two.

I. Introduction

The systematic differences between the high energy total cross sections and elastic scatterings of $\bar{K}^-p$ or $\bar{p}p$ and their counterparts under crossing, K^+p or pp , have long been remarked and discussed. For the former systems of type $\bar{X}p$, one has larger total cross sections σ_T which decrease slowly with increasing energy and the diffraction peaks are sharper. For the latter of type Xp , total cross sections are smaller and almost energy independent. The differences are very large for $\bar{p}p$ versus pp and quite appreciable for $\bar{K}^-p$ versus K^+p . Although smaller and less striking, similar differences exist for π^-p versus π^+p .

Observation and physical intuition suggest that high energy elastic scattering is predominantly a diffractive process, being the shadow effect in the elastic channel of the many inelastic processes which make up some 80% of the total cross section. It is therefore natural to look to systematic differences in the inelastic processes for an explanation of the above features. This idea is particularly natural for $\bar{p}p$ versus pp , where the baryon annihilation processes $\bar{p}p \rightarrow \text{mesons}$, with their high cross section, are obvious candidates to explain the larger σ_T and, consequently, the sharper diffraction peak of $\bar{p}p$. In fact, at the highest energies where it was measured, the baryon annihilation cross section is found to be approximately equal to the difference $\delta\sigma_p = \sigma_T(\bar{p}p) - \sigma_T(pp)$; the former cross section is 23.6 ± 3.4 mb at beam momentum $q_L = 7$ GeV/c [1], the latter difference being about 20 mb [2].

Similarly, we shall attempt to relate the differences between $\bar{K}^-p$ and K^+p to the "hypercharge annihilation" processes $\bar{K}^-p \rightarrow \text{hypercharge zero particles}$ (mainly one $Y = 0$ hyperon and a number of $Y = 0$ mesons). It is known that the total cross section for $\bar{K}^-p \rightarrow \Lambda + \text{pions}$ is 2.3 ± 0.2 mb at 10 GeV/c [3]. This should be increased by the unknown cross section for $\Sigma^\pm$ production and might amount to at least half of the difference

$$\delta\sigma_K = \sigma_T(\bar{K}^-p) - \sigma_T(K^+p) \simeq 5.1 \text{ mb [2] at the same energy. It}$$

is doubtful that it could amount to the totality of this difference, because the mechanism responsible for the $\delta\sigma_\pi = \sigma_\pi(\pi^-p) - \sigma_\pi(\pi^+p)$ difference is likely to have an analogue for Kp and to give to $\delta\sigma_K$ a contribution of order $\delta\sigma_\pi$ which is 1.7 mb at 10 GeV/c [2].

To explain the latter mechanism, it is more difficult to propose a single class of inelastic processes comparable to the two forms of annihilation just considered. The concept of charge annihilation defined as $\bar{X}p \rightarrow \text{neutrals}$ cannot be adopted because it is not invariant for isospin rotations (remember e.g. that π^-p and π^+n are identical up to an isospin rotation). What is needed is a concept expressing the general idea that more inelastic processes of charge exchange type are accessible to π^-p than to π^+p , more to π^0p than to π^+p , more to K^0p than to K^+p , etc. Our basic assumption is that such additional charge exchange processes contribute to the $\bar{X}p$ versus Xp cross section differences; in particular, they are responsible for most of the π^+p difference. For convenience we give to these extra charge exchange channels the name of "charge annihilation". We stress, however, that the concept is not sharply defined and that the terminology is purely ad hoc.

The approach just mentioned has been first discussed for the $\bar{p}p$ and pp systems by Kokkedee [4], and then, in considerable detail with inclusion of spin effects, by one of us for πp [5,6]. The present paper treats it for Kp and gives a more complete discussion of $\bar{p}p$, pp . The main outcome is the qualitative prediction of cross section and range of the annihilation processes (mainly baryon annihilation for $\bar{p}p$, hypercharge + charge annihilation for $\bar{K}^-p$) from the available high energy data for σ_π and $d\sigma_{el}/dt$. We shall find that for $\bar{p}p$ the range of annihilation is roughly equal to the range of elastic scattering. The annihilation range has again about the same numerical value for $\bar{K}^-p$, but in this case the elastic range is smaller by approximately a factor 2. We further predict the decrease of total annihilation cross sections at higher energies.

Before describing the principle of our method, we briefly contrast our approach with the more recent proposal of Harari that the differences of $\bar{X}p$ versus Xp properties at high energy are entirely describable as contributions from s-channel resonances [7]. This hypothesis has been tested on the π^+p system, but the result depends on how resonances are parametrized, as illustrated e.g. in recent papers by Dance and Shaw [8] on the one hand and by Harari and Zarmi [9] on the other. While the Harari scheme is not in logical contradiction with ours, it seems perhaps less natural when one realizes that it runs into difficulties for $\bar{p}p$ versus pp , i.e. for the very case where the differences to be explained are largest (for these difficulties, see Rosner [10] and sections II.6 and VII.6 of ref. 9). Furthermore, our approach can be tested experimentally at high energy by measuring s-channel annihilation processes, whereas Harari's requires knowledge of the high mass resonance parameters, which are bound to remain largely inaccessible to experiment for masses $\gtrsim 3$ GeV. It should be stressed that our scheme of "annihilation duality" has essentially the same consequences as Harari's duality for exchange degeneracy of Regge trajectories and residues, as noted by Chiu and Finkelstein [11]. In fact, Henzi's treatment of the πp system [5,6] was the first to give exchange degeneracy of both trajectories and residues for ρ and P' exchange.

The central idea of our scheme is then that the main non-asymptotic part of elastic scattering (in Regge language this is its non-Pomeranchuk part) is the shadow effect of the generalized annihilation processes in the s-channel. Its implementation relies on the overlap function concept, a function $G(b,s)$ which describes the contribution of all inelastic channels to the cross section for given impact parameter b . It is additive in the different inelastic channels and contains therefore, in the case of $\bar{X}p$, a term δG_B contributed by the annihilation channels mentioned above (we call these also B-type channels). The other, or A-type channels are supposed to be very similar for Xp and $\bar{X}p$ and to contribute a term G_A . For Xp , the data allow one to put G_A equal to the

common infinite energy limit G_∞ of G . For $\bar{X}p$, G_A is written as $G_\infty - \delta G_A$ where δG_A is a modification of G_A attributed to the reaction effect on A-type channels produced by the presence of the B-type ones. Using a simple parametrization to analyse the data on σ_T and elastic scattering, we extract from them the quantity

$$\delta \zeta(b,s) = \delta \zeta_B(b,s) - \delta \zeta_A(b,s).$$

We then propose physical models which allow one to separate δG into δG_A , δG_B . This separation turns out to be reasonably insensitive to the model used. Obtaining $\delta G_B(b,s)$ in this way, we have information on total cross section and range of the annihilation channels, both of which can in principle be confronted with experimental data. The consistency of our approach, especially of the assumption that A-type channels are similar for Xp and $\bar{X}p$, is supported a posteriori by the fact that δG_A turns out to be much smaller than δG_B for most impact parameters b .

The details of our method will be found in the next section, and the quantitative analysis of the $K^\pm p$, $\bar{p}p$ and pp data is carried out in section III. The last section is devoted to concluding remarks. The models used for the separation of δG into δG_A , δG_B are discussed in an appendix.

II. Assumptions and Method

First, we recall the formalism and assumptions for the overlap function analysis. Spins being neglected ^{*}), the differential cross section is given in terms of the scattering amplitude (suitably normalized) by the relation

$$d\sigma_{el}/dt = |F(s,t)|^2. \quad (1)$$

At high energies, it is convenient to analyse $F(s,t)$ in terms of the impact parameter expansion

$$F(s,t) = 2\pi^{1/2} \hbar \int_0^\infty h(b,s) J_0[b(-t)^{1/2}] b db \quad (2)$$

with the weight function $h(b,s)$ given by the inverse relation

$$h(b,s) = (4\pi^{1/2} \hbar)^{-1} \int_0^\infty F(s,t) J_0[b(-t)^{1/2}] d(-t). \quad (3)$$

As shown by one of us [5], these formulae are valid up to corrections of order $s^{-1+\epsilon}$, ϵ positive and arbitrarily small. They are sufficient for our purpose. We introduce $\hbar$ into the formulae so that we may work in mixed units with cross sections σ measured in millibarns, t in $(\text{GeV}/c)^2$, b in $(\text{GeV}/c)^{-1}$. Numerically,

$$\hbar^2 = 0.389 (\text{GeV}/c)^2 \text{ mb}. \quad (4)$$

At high energies, unitarity is approximately given by the relation

$$\text{Im } h(b,s) = \frac{1}{4} G(b,s) + |h(b,s)|^2 \quad (5)$$

^{*}) A full spin treatment has been given by Henzi for the πp system [5,6]

where $G(b,s)$, the overlap function, determines the elastic shadow scattering due to all open inelastic channels. It is a sum of positive terms, one for each of these channels.

The formulae for the cross sections of interest are then (in the usual notation)

$$\sigma_T = 4\pi^{1/2} \hbar \operatorname{Im} F(s,0) = 8\pi \hbar^2 \int_0^\infty \operatorname{Im} h(b) b db \quad (6)$$

$$\sigma_{inel} = \sigma_T - \sigma_{el} = 2\pi \hbar^2 \int_0^\infty G(b) b db \quad (7)$$

$$d\sigma_{el}/dt = 4\pi \hbar^2 \left| \int_0^\infty h(b) J_0[b(-t)^{1/2}] b db \right|^2 \quad (8)$$

$$(d\sigma_{el}/dt)_{t=0} = (16\pi \hbar^2)^{-1} \sigma_T^2 (1+\alpha^2) = (\sigma_T^2/19.6)(1+\alpha^2) \quad (9)$$

where $\alpha = \operatorname{Re} F(s,0) / \operatorname{Im} F(s,0)$. (10)

The experimental information to which F or h are to be fitted comprises

(1) σ_T

(2) σ_{el} / σ_T

(3) α or equivalently $(d\sigma_{el}/dt)_{t=0}$

(4) The shape of $d\sigma_e/dt$ for moderate t , say $|t| \lesssim 1.2 (\text{GeV}/c)^2$.

The shape function is commonly parametrized as an exponential

$$[d\sigma_e(s,t)/dt] / [d\sigma_e(s,0)/dt] = \exp(\beta t + \gamma t^2 + \dots) \quad (11)$$

where the linear term in the exponent is the strongest feature for moderate t , yielding drops of about three orders of magnitude out to $-t \simeq 1.2 (\text{GeV}/c)^2$.

(5) More recently, data has extended to wider angles (larger $|t|$). The general feature has been a flattening out of the exponential drop, and for some processes the occurrence of secondary diffraction maxima.

The models to be employed in our analysis belong to the general class of optical models in which the high energy scattering amplitude F is assumed to approach a fixed pure imaginary function of t as the beam energy tends to infinity. We are interested in the main trends of the elastic scattering parameters as a function of energy as they approach the infinite energy limits indicated by their extrapolation. We want to interpret the deviation from the asymptotic limits in terms of special classes of inelastic channels, the annihilation channels already mentioned in the introduction. By use of the overlap function we shall correlate some properties of these channels with the observed σ_T and elastic scattering. We now discuss in succession the various steps of the analysis to be applied in the next section.

a) Infinite energy limit

For the infinite energy limit of the overlap function, taken to be the same for X_p and $\bar{X}_p$, we shall mostly adopt the simple Gaussian shape

$$G_\infty(b) = P \exp(-b^2/4A) \quad (12)$$

corresponding to a pure imaginary weight function h_∞ given by

$$-i h_\infty = \text{Im } h_\infty = \frac{1}{2} [1 - (1 - G_\infty)^{1/2}] \quad (13)$$

The corresponding S-matrix is given by

$$S_\infty(b) = 1 - 2 \text{Im } h_\infty = (1 - G_\infty)^{1/2} \quad (14)$$

This model [12] predicts values for the ratio σ_{el}/σ_T in accord with experiment and gives an almost exponential drop of $d\sigma_{el}/dt$, with small positive values for the quadratic coefficient γ of eq. (11). It represents the "scatterer" as being very smooth with gradual edges and consequently does not predict any secondary diffraction structure. The model has been applied to a number of experimental situations [4,5,6,13].

The accumulation of elastic scattering data at wider angles in which the sharp experimental drop was seen to abate or even, in some cases, to give place to secondary diffraction minima, stimulated Chou and Yang to introduce their droplet model [14]. This achieves a remarkable correlation with the electromagnetic form factors. All that concerns us here is that a more "bell-shaped" weight function is invoked leading to secondary diffraction maxima even at infinite energy. We shall briefly discuss the implications for our analysis of $\bar{p}p$ and pp , using the explicit formulae of Durand and Lipes [15]. These give for the overlap function

$$G_\infty(b) = 1 - \exp[-2\lambda \phi(\mu b)] \quad (15)$$

with

$$\phi(x) = \frac{1}{8} x^3 K_3(x). \quad (16)$$

We note nevertheless that the evidence for the existence of secondary diffraction maxima at infinite energy is weak. As illustrated by Czyz (see section 3 and Fig. 9 of ref. 16), the observed structure in pp elastic scattering at 19.3 GeV need not be interpreted as the first indication of such a maximum and can equally well be understood as a finite energy effect.

b) Two contributions to δG

We now turn to the question of the way G approaches G_∞ for Xp as against $\bar{X}p$ scattering. As already stated, we assume that the overlap function for Xp scattering [reaction (0)] approaches the asymptotic limit G_∞ more rapidly than for $\bar{X}p$ [reaction (1)], and we approximate $G^{(0)}$ by G_∞ at beam momentum $q_L \gtrsim 5$ GeV/c. We thus write

$$G^{(0)} = G_\infty \quad (17)$$

In contrast, $\bar{X}p$ has an important energy dependent term δG , i.e.

$$G^{(0)} = G_\infty + \delta G, \quad \delta G \rightarrow 0 \text{ for } s \rightarrow \infty \quad (18)$$

As mentioned before, this is essentially the same concept as employed in Harari's duality for $X = K^+$ and p , the only difference being that the latter assumes constancy of $\text{Im}h^{(0)}$, which differs from (17) by a term $[\text{Re}h^{(0)}]^2$, see eq. (5). At $b=0$, Harari's duality means constancy of $\sigma_T(Xp)$ whereas our model means constancy of $\sigma_{inel}(Xp)$.

We note at this stage that both our scheme and Harari's duality will have to be revised if $\sigma_T(Xp)$ is found to increase at high energy. For us, this would simply mean that the cross section for A-type processes would have an hitherto undetected increase with

energy, which could for example result from opening up of new channels. Speculations in this direction are encouraged by recent Serpukhov measurements of σ_T for $\bar{p}p$, $\bar{K}^-p$ and π^-p which indicate that the decrease of these cross sections for increasing energy may be less than expected [17]. To save the Pomeranchuk theorem, a slow increase of $\sigma_T(pp, Kp, \pi^+p)$ may therefore become desirable. For our present analysis of the data (see next section), use of the Serpukhov results for $\bar{X}p$ is impossible without speculative assumptions on Xp reactions, and we shall restrict ourselves to the energy region where both Xp and $\bar{X}p$ are known.

We now list our assumptions on $\delta G = G^{(1)} - G_\infty$:

- i) δG is attributable to inelastic channels accessible to $\bar{X}p$ which have no counterpart in Xp (B-type processes). They have cross sections tending to zero for increasing energy. These channels contribute a term δG_B to the overlap function of $\bar{X}p$.
- ii) The other, and predominant inelastic channels (A-type processes) are similar in Xp and $\bar{X}p$. While their contribution to $G^{(0)}$ is approximately asymptotic, eq. (17), their contribution to $G^{(1)}$ is reduced by a reaction effect due to the presence of the B-type processes. We write it in the form $G_\infty - \delta G_A$, so that

$$\delta G = \delta G_B - \delta G_A \quad (19)$$

The necessity of this reaction effect is clear when G_∞ is close to the unitary limit. We expect $0 < \delta G_A < \delta G_B$.

For $\bar{p}p$, the main candidates we have in mind for B-type processes are the baryon annihilation channels $\bar{p}p \rightarrow \text{mesons}$. We could also consider as B-type processes the channels $\bar{p}p \rightarrow \text{hypercharge zero particles ("hypercharge annihilation")}$, or charge exchange channels missing in pp ("charge annihilation"), but they are much less abundant and their contribution to $\mathcal{G}_B$ and $\mathcal{G}_A$ is therefore unimportant at our level of approximation. We assume that all other inelastic channels important at high energy are of A-type, i.e. exist for both pp and $\bar{p}p$, except for the replacement of a nucleon by an antinucleon. They are presumably of form $pp \rightarrow 2N + \text{pions}$, $\bar{p}p \rightarrow \bar{N}N + \text{pions}$ (the latter still containing charge annihilation).

For $\bar{K}^-p$, B-type processes are expected to be mainly hypercharge and charge annihilation, i.e. $\bar{K}^-p \rightarrow \text{hyperons} + \text{non-strange mesons}$ and charge exchange channels open to $\bar{K}^-p$ but missing in K^+p . The A-type processes would be mainly of type $K^+p \rightarrow K + N + \text{non-strange mesons}$, $\bar{K}^-p \rightarrow \bar{K} + N + \text{non-strange mesons}$ (the latter also containing charge annihilation). Regarding hypercharge annihilation $\bar{K}^-p \rightarrow \text{hyperon} + \text{non-strange mesons}$, attention should be paid to the fact that the hyperon could be a Y^* resonance decaying into $\bar{K} + N + \text{possibly pions}$. It would then be detectable through the presence of a backward $\bar{K}$.

For π^-p we have only what we call charge annihilation. We note here in passing that the cross section for $\pi^-p \rightarrow \text{neutron} + \pi^0$'s (which would be charge annihilation in the literal sense) is quite small. This can be deduced from measurements of the total neutral cross section σ_n in π^-p interactions in the interval $1.4 < q_L < 4.0 \text{ GeV}/c$ [18]. It is found that σ_n drops rapidly with increasing beam momentum and is $\simeq 1.1 \text{ mb}$ at $q_L = 4 \text{ GeV}/c$, i.e. much too small to account for $\sigma_T(\pi^-p) - \sigma_T(\pi^+p)$.

Isospin invariance implies relations among all πN reactions, and similarly among XN or $\bar{X}N$ ones ($X = K$ or N). They are best taken into account in our formalism by defining G_I as the overlap

function for an initial state XN or $\bar{X}N$ of given isospin I in the s -channel. For πN one then has

$$\text{for } \pi^+ p, G^{(0)} = G_{3/2}; \quad \text{for } \pi^- p, G^{(1)} = \frac{1}{3} G_{3/2} + \frac{2}{3} G_{1/2}$$

and charge annihilation refers to the additional charge exchange channels which are open in $G_{1/2}$, i.e. when the incoming isospins are not parallel. More generally, for XN or $\bar{X}N$ ($X = K$ or N), charge annihilation will contribute to the overlap function whenever the incident isospins are not parallel, i.e. to G_I for $I < I_N + I_X$ but not to G_I for $I = I_N + I_X$.

While δG can be determined from elastic scattering data [see d) below], the separation of δG_A and δG_B requires a model. Given such a model, δG_B can be calculated from the elastic data and allows one to predict the total cross section of the annihilation processes as well as their distribution in impact parameter, i.e. their spatial range. Such predictions can be tested experimentally. Another application would be to compare δG_B with the overlap function contribution of B-type processes as could be calculated from detailed models, e.g. the multi-Regge model.

c) Models for $\delta G_A, \delta G_B$

The crucial question is of course whether reasonable models exist for δG_A and δG_B , leading to a sufficiently unique determination of these functions from δG . We have considered three such models; as shown in the next section the analysis of the data gives $\delta G_A, \delta G_B$ -values which agree well with each other.

1) The first or Gaussian model was introduced by one of us for the discussion of the πp system [5,6]. The shape of δG in b^2 , as deduced from the precise πp data, suggested the simple parametrization

$$\delta\zeta = (1 - \zeta_\infty)^{1/2} [Q \exp(-b^2/4B) - Q' \exp(-b^2/4B')]. \quad (20)$$

We find that such a parametrization applies also very well to $\bar{K}p$ and $\bar{p}p$. It leads to the natural splitting

$$\delta\zeta_B = (1 - \zeta_\infty)^{1/2} Q \exp(-b^2/4B), \quad \delta\zeta_A = (1 - \zeta_\infty)^{1/2} Q' \exp(-b^2/4B'). \quad (21)$$

The physical interpretation of these formulae is simple. The B-type channels have a Gaussian distribution in b , corrected for absorption due to the much stronger A-type channels. The absorption correction is given the familiar form of multiplication by the S-matrix element $S(b)$ for which the infinite energy value (14) is taken. The reaction effect δG_A is given the same shape. Since it cannot extend beyond the ranges of either A- or B-type processes, we require

$$B' \leq A, \quad B' \leq B. \quad (22)$$

It is also clear that we must require positive $\delta G_A, \delta G_B$, i.e.

$$Q > 0, \quad Q' > 0. \quad (23)$$

These conditions will all be satisfied in our fits.

Our two other models are a little more theoretical in nature. They express δG_A and δG_B in terms of the hypothetical B-type overlap function δg which would obtain if the A-type processes were switched off. This is done in two ways, one using a K-matrix formulation and the other an eikonal description of a continuous medium. Both models, which turn out to give very similar results in the fits of the next section, are outlined in some detail in the Appendix. Here we only give the final formulae, where S_∞

is given by eq. (14)

$$S_{\infty} = (1 - G_{\infty})^{1/2}.$$

2) K-Matrix model

$$\delta G = \frac{1}{4} (1 - S_{\infty})^{-2} \cdot S_{\infty} \cdot G_{\infty}^2 \cdot \delta g \quad (24)$$

$$\delta G_B = \frac{1}{4} (1 - S_{\infty})^{-2} \cdot G_{\infty}^2 \cdot \delta g = \delta G / S_{\infty} \quad (25)$$

$$\delta G_A = \frac{1}{4} (1 - S_{\infty})^{-1} \cdot G_{\infty}^2 \cdot \delta g = (1 - S_{\infty}) \cdot \delta G / S_{\infty} \quad (26)$$

3) Eikonal model

$$\delta G = S_{\infty}^2 \cdot \delta g \quad (27)$$

$$\delta G_B = -(\ln S_{\infty}^2)^{-1} \cdot G_{\infty} \cdot \delta g = (-S_{\infty}^2 \ln S_{\infty}^2)^{-1} G_{\infty} \cdot \delta G_{(28)}$$

$$\delta \zeta_A = - \left(S_\infty^2 + \frac{G_\infty}{\ln S_\infty^2} \right) \delta g = - \left(1 + \frac{G_\infty}{S_\infty^2 \ln S_\infty^2} \right) \delta \zeta \quad (29)$$

For both models, we have at large b^2

$$\delta \zeta_B \simeq \delta \zeta, \quad \delta \zeta_A \simeq \frac{1}{2} G_\infty \cdot \delta \zeta \quad (30)$$

so that the range conditions expressed before eq. (22) are satisfied. Positivity of δG_A , δG_B results from the fact that the empirical $\delta G(b)$ will be found positive.

d) Parametrization and practical procedure

The first step of our procedure will consist in determining the two free parameters of $G_\infty(b)$ by estimating from the data the infinite energy limits of σ_T and of the slope of $d\sigma_{el}/dt$ at $t=0$, which is expressed by the parameter β of eq. (11). Our model assumes that the amplitudes $F^{(0)}$, $F^{(1)}$ for XP , $\bar{X}p$ have a common pure imaginary limit

$$F_\infty(t) = i \operatorname{Im} F_\infty(t). \quad (31)$$

We then analyse the energy dependent effects within the following simple approximation (a more precise analysis was possible for πp , as shown in refs. 5,6). For Xp , the amplitude F is given the form

$$F^{(0)} = F_\infty(t) + q_L^{-\nu} \cdot f^{(0)}(t) \quad (32)$$

with q_L the laboratory momentum, ν a positive constant and $f^{(0)}$ a real function of t only. Up to corrections of order $q_L^{-2\nu}$, the corresponding form of the overlap function $G^{(0)}$ calculated by means of eqs. (3) and (5) reduces to the infinite energy function G_∞ as indicated in eq. (17).

For reaction $\bar{X}p$, the q_L dependence of the amplitude is written in the form

$$F^{(1)} = i \left[\text{Im } F_\infty(t) + c q_L^{-\nu} \cdot Y(t) \right] + q_L^{-\nu} \cdot f^{(1)}(t) \quad (33)$$

with c a real constant, and Y , $f^{(1)}$ real functions of t only, Y satisfying in addition

$$Y(0) = 1 \quad (34)$$

ν is the same as in eq. (32). Up to corrections of order $q_L^{-2\nu}$, the corresponding overlap function $G^{(1)}$ is obtained as follows. From eq. (5), neglecting the contribution of $\text{Re } h^{(1)}$, one gets

$$\begin{aligned} G^{(1)} &= G_\infty + \delta G = 4 \text{Im } h^{(1)} - 4 (\text{Im } h^{(1)})^2 \\ &\simeq G_\infty + 4 (\text{Im } h^{(1)} - h_\infty) (1 - 2 h_\infty) \\ &= G_\infty + 4 (\text{Im } h^{(1)} - h_\infty) (1 - G_\infty)^{1/2} \end{aligned}$$

Eq. (14) has been used in the last step. From (33) and (3) we now get

$$\delta G = 2c (\pi^{1/2} k)^{-1} (1 - G_\infty)^{1/2} \cdot q_L^{-\nu} \cdot Z(b) \quad (35)$$

$$Z(b) = \frac{1}{2} \int_0^\infty Y(t) J_0 [b(-t)^{1/2}] d(-t) \quad (36)$$

Eqs. (32) and (33) give for the total cross sections, up to terms in $q_L^{-2\nu}$,

$$\sigma_T^{(0)} = \sigma_T^{(\infty)}, \quad \sigma_T^{(1)} = \sigma_T^{(\infty)} + \delta\sigma \quad (37)$$

$$\delta\sigma = 4\pi^{1/2} \hbar c q_L^{-\nu} \quad (38)$$

The high energy data for σ_T can be fitted with these approximate formulae and determine c and ν . For the elastic scattering our formulae give

$$d\sigma_{el}^{(0)}/dt = [F_\infty(t)]^2 \quad (39)$$

$$d\sigma_{el}^{(1)}/dt = [F_\infty(t) + 2c q_L^{-\nu} Y(t)]^2 \quad (40)$$

where the contributions of the real parts of F to $d\sigma_{el}/dt$, which are of order $q_L^{-2\nu}$, have been neglected. We note at this stage that analyticity and crossing symmetry allow one to derive $f^{(0)}$ and $f^{(1)}$ from Y

$$f^{(0)} = -cY/\sin \pi\nu, \quad f^{(1)} = -f^{(0)} \cos \pi\nu \quad (41)$$

We shall obtain for γ values of order 0.6, so that

$$f^{(0)} \simeq -1.05 c Y, \quad f^{(1)} \simeq -0.32 c Y$$

The latter equation shows that the real part contribution is much smaller in $d\sigma_{el}^{(1)}/dt$ than in $d\sigma_{el}^{(0)}/dt$. We therefore choose to derive Y from $d\sigma_{el}^{(1)}/dt$ by neglecting $f^{(1)}$, i.e. by means of the simple formula

$$Y(t) = \frac{[d\sigma_{el}^{(1)}(s,t)/dt]^{1/2} - [d\sigma_{el}^{(\infty)}(t)/dt]^{1/2}}{[d\sigma_{el}^{(1)}(s,0)/dt]^{1/2} - [d\sigma_{el}^{(\infty)}(0)/dt]^{1/2}} \quad (42)$$

Since the data do not show any significant energy dependence of this expression, we shall combine them for reactions at various energies.

Finally, it appears that the data for eq. (42) are easily fitted to

$$Y(t) = (1+H) \exp(Bt) - H \exp(B't) \quad (43)$$

with H, B, B' positive constants. Through eqs. (35), (36) one then obtains for δG the form (20) discussed previously, more precisely

$$\delta G = (4\pi k^2)^{-1} \Phi(b) \cdot \delta \sigma \quad (44)$$

$$\Phi(b) = (1 - \zeta_{\infty})^{1/2} \left[B^{-1} (1+H) \exp(-b^2/(4B)) - B'^{-1} H \exp(-b^2/(4B')) \right]$$

The split of δG into $\delta G_A, \delta G_B$ proceeds as discussed under c) above by means of the models there mentioned. We shall write

$$\delta \zeta_{A,B} = (4\pi k^2)^{-1} J_{A,B}(b) \cdot \delta \sigma \quad (45)$$

$$\sigma_{inel}^{A,B} = J_{A,B} \cdot \delta \sigma = \frac{1}{4} \left[\int_0^\infty J_{A,B}(b) db^2 \right] \cdot \delta \sigma \quad (46)$$

J_B expresses the physically interesting ratio of σ_{inel}^B , the total cross section for the B-type processes, to the total cross section difference $\delta \sigma = \sigma_T(\bar{X}p) - \sigma_T(Xp)$. As mentioned before, consistency of the model requires

$$0 < \delta \zeta_A \ll \delta \zeta_B, \quad 0 < \sigma_{inel}^A \ll \sigma_{inel}^B, \quad 0 < J_A \ll J_B.$$

Since σ_{el} is always small compared to σ_{inel} , $\delta \sigma$ will be approximately equal to $\sigma_{inel}^{(1)} - \sigma_{inel}^{(0)}$, and from eq. (7) this is the same as $\sigma_{inel}^B - \sigma_{inel}^A = (J_B - J_A) \delta \sigma$. We therefore expect J_B to be close to 1 if our model is to be consistent.

If one wishes to compare our parametrization to the conventional Regge pole analysis, one sees that we simplify on the following points: exchange degeneracy of secondary trajectories, neglect of shrinking (i.e. of the trajectory slopes), and neglect of spin terms. The two latter effects were included in the more precise analysis of πp scattering [5,6]. The available data do not warrant their study for the pp , $\bar{p}p$ and Kp systems.

III. Fits to the Data

We now turn to the task of determining $\delta G_{A,B}$ by the method outlined in the previous section. To recapitulate, the steps are as follows:

- i) Fix parameters P and A of $G_{\infty} = P \exp(-b^2/4A)$ from extrapolations of σ_T and β , the logarithmic slope of $d\sigma_{el}/dt$ at $t=0$. For pp , we indicate how the results are modified by using the Durand-Lipes form (15), which uses the parameters λ, μ instead of P and A .
- ii) Compute $d\sigma_{el}^{(\infty)}/dt$ from G_{∞} , determine $Y(t)$ from the data using (42) and fit it to the expression (43). This gives H , B and B' .
- iii) Determine c and ν from the σ_T data, by fitting $\delta\sigma$, eq. (38).
- iv) Calculate Φ and δG from (44). Split δG into $\delta G_B - \delta G_A$, i.e. Φ into $\Phi_B - \Phi_A$ using the three models described in section IIc, eqs. (21), (25), (26), (28), (29). Calculate $J_{A,B}$ and $\sigma_{inel}^{A,B}$ from (46).
- v) Compare and discuss the results.

We now treat in succession Kp and $pp, \bar{p}p$.

1) Kp Scattering

The K^+p total cross sections from 6 to 20 GeV/c are remarkably constant [2], fluctuating about the value 17.2 mb, so we assume

$$\sigma_T^{(\infty)} = 17.2 \text{ mb}$$

In seeking to determine our second asymptotic parameter by using the experimental slope parameter β , we encounter a practical difficulty. The experimental values for the K^+p differential cross section [19], when extrapolated to $t=0$, imply values of $\text{Re}F$ far in excess of what would appear reasonable from forward dispersion relations [20], even allowing for the quoted uncertainties in the absolute normalisation. It appears that the early analysis of ref. 19 would now be viewed as inadequate in such a way as to render too large the deduced values of $d\sigma_{el}/dt$ at small $|t|$ [21].

Our standpoint is therefore to fix the forward point from the optical theorem. We take the K^+p total cross sections from ref. 22 and these are listed along with other cross section data in Table I. We take $\alpha = (\text{Re}F/\text{Im}F)_{t=0}$ from forward dispersion relation calculations [20]. The values which we adopt and the resulting $(d\sigma_{el}/dt)_{t=0}$ are listed in Table II. We now determine the β parameters by fitting the published $d\sigma_{el}/dt$ data in the following two ways: (i) Constrain the forward point and refit (this implies a belief in important systematic errors at small $|t|$); (ii) Fit without constraining the forward point (this corresponds to the assumption of a t -independent normalisation error).

The resulting β values are shown in Fig. 1 (Method (i) with full-line error bars, Method (ii) with dotted error bars). The points are plotted against q_o/q_L with q_L the incident beam momentum and $q_o = 10 \text{ GeV}/c$. Also shown on the same figure are the corresponding points for $\bar{K}^-p$, in this case plotted against $[\sigma_T^{(1)}(s) - \sigma_T^{(\infty)}]/\sigma_T^{(1)}(s)$. (The abscissa against which the $\bar{K}^-p$ slope is plotted is chosen to give a linear dependence of β on a simple Regge model of the non-asymptotic effects as postulated in eq. (40); for K^+p , we simply took the abscissa proportional to q_L^{-1} . The scales are unconnected. From the generalized Pomanchuk theorem, one expects the two sets of points to lie on approximately straight lines of distinct slope but common intercept at $q_L = \infty$). Also shown on the side of Fig. 1 are the corresponding values of the opacity parameter P appearing in G_∞ , eq. (12), when $\sigma_T(\infty) = 17.2 \text{ mb}$ is assumed.

Inspection of Fig. 1 indicates a value of β in the range 5-7 (GeV/c)⁻² corresponding to P between 1.0 and 0.8. We then compare in Fig. 2 the predicted $d\sigma_{\ell}^{(\infty)}/dt$ with K^+p data points. The effect of the real part of the amplitude is seen to be still appreciable [a feature which is perhaps not surprising on the basis of the estimate given after eq. (41)]. It is nevertheless possible to use Fig. 2 for reducing somewhat the uncertainty on P , and it seems that the range $P = 0.85$ to 0.95 is acceptable in the sense of not postulating excessively large real parts or large errors. (We are now appealing to the height of $d\sigma_{\ell}^{(\infty)}/dt$ as well as its slope). $P = 0.85$ gives about the right slope but requires for all t a large normalization error or a large real part; $P = 0.95$ requires an error or a real part which are large at small $|t|$. We can narrow the possibilities further by appeal to the $\bar{K}^-p$ data ^{*)}. A number of alternative fits to $G^{(1)}(b^2)$ which we have made (assuming $\text{Re}F^{(1)}$ negligible) agree that $G^{(1)}(b^2=0)$ is of the order of 0.90 or less. It seems reasonable to assume $G^{(1)} \geq G_{\infty}$ at all b^2 , in particular at $b^2=0$. We therefore finally conclude that the reasonable range for $P = G_{\infty}(0)$ is 0.85 to 0.90. $P = 0.90$ implies a pronounced $K^+p/\bar{K}^-p$ "cross-over" at $|t| \sim 0.2$; $P = 0.85$ implies that there is no clear cross-over (recall that this figure required large systematic normalization errors for K^+p). We shall now proceed to work with the two values $P = 0.90$ and $P = 0.85$. The corresponding values of the parameter A of (12) for $\sigma_T^{(\infty)} = 17.2$ mb are 3.32 and 3.53 (GeV/c)⁻².

To determine c and ν , eqs. (37,38), we fitted the $\bar{K}^-p$ total cross section data up to 20 GeV/c [2] to the form

$$\sigma_T^{(1)} = \sigma_T^{(\infty)} + 21.5 q_L^{-0.61} = 17.2 + 21.5 q_L^{-0.61}$$

*) $\bar{K}^-p$ differential cross section data used refers to beam momenta 7.2 and 9.0 GeV/c [19], 10.12 GeV/c [23], 11.88 and 15.91 GeV/c [24]. Only data for $|t| \leq 0.7$ (GeV/c)² was used.

with σ_T in mb and q_L in GeV/c. This fit does not describe well the Serpukhov data for $\sigma_T(\bar{K}^-p)$ [17]; we have explained in section IIb why we prefer to ignore them until $\sigma_T(K^+p)$ has been measured at comparable energies.

We next compute $Y(t)$ from the $\bar{K}^-p$ elastic data, using eq. (42). The values of $Y(t)$ for the case $P = 0.90$ are shown in Fig. 3a. These values must be fitted to formula (43). They do not precisely determine the three separate parameters H , B and B' but admit a range of possibilities with H , B and B' tightly correlated. We list in Table III the parameters for three representative solutions more or less spanning the range of possibilities and plot these fits along with the data in Fig. 3a. The experimental points for $P = 0.85$ are shown in Fig. 3b together with a representative solution for this case whose parameters are also listed in Table III. Inspection of the table reveals an important qualitative conclusion: B , the range of the B-type processes has approximately twice the value of A , the range of the asymptotic A-type processes. This result essentially derives from the upward curvature of the $\bar{K}^-p$ differential cross section at small $|t|$ for which the data is compared with $d\sigma_a^{(\infty)}/dt$ in Fig. 4. This effect is absent in K^+p , see Fig. 2, and our model therefore allocates it to $Y(t)$. It concerns the small $|t|$ range and hence corresponds to a long range B in b^2 .

Given the parameters H , B and B' , we can go on to calculate δG by (44) and separate it into A- and B-type contributions, $\delta G_{A,B} = (\delta\sigma/4\pi\hbar^2)\bar{\Phi}_{A,B}$, according to our three alternative prescriptions of eqs. (21), (25,26) and (28,29). Further, by a simple quadrature, we calculate the parameters $J_{A,B}$, eq. (46), which determine the A- and B-type contributions to the total inelastic cross sections.

The resulting forms for $\bar{\Phi}_{A,B}$ according to the three recipes are plotted in Fig. 5a for fit II, and the forms given by method 2 for fits I to IV in Fig. 5b. The results for J_B are listed in Table III. Methods 2 and 3 give a rather constant pre-

diction for the four fits listed. It satisfies $J_A \ll J_B$ as required for the consistency of our model. Method 1 is naturally more sensitive to the fit parameters since the separation is made directly in terms of them. It gives larger J_A , especially when H is taken large, and solution I is quite unsatisfactory for this method. For the purpose of discussion, we take the average

$$\langle J_B \rangle \simeq 1.1 \pm 0.25.$$

On our model, the quantity σ_{inel}^B is to be that part of the $\bar{K}^-p$ inelastic cross section attributable to B-type processes, i.e. hyperon and charge annihilation. Our predictions using the above average value $\langle J_B \rangle$ are listed in Table IV together with the relevant data. At 6 and 10 GeV/c there are measurements of $\sigma(\bar{K}^-p \rightarrow \Lambda + \pi's)$ [25,3]. The total hypercharge annihilation cross section $\sigma_Y(\bar{K}^-p)$ will be larger due to $\Sigma^\pm$ production and to production of Y^* resonances decaying into N , K and possibly pions. We have no way of measuring the charge annihilation cross section for $\bar{K}^-p$, but we can assume it to be comparable to the differences

$$\delta_1 = \sigma_T(\bar{K}^-p) - \sigma_T(\bar{K}^-n), \quad \delta_2 = \sigma_T(\pi^-p) - \sigma_T(\pi^+p)$$

This means that we predict

$$\begin{aligned} \sigma_{inel}^B(\bar{K}^-p) &\simeq \sigma_Y(\bar{K}^-p) + \delta_1 \simeq \\ &\simeq \sigma_Y(\bar{K}^-p) + \delta_2 \gtrsim \sigma(\bar{K}^-p \rightarrow \Lambda + \pi's) + \delta_2 \end{aligned}$$

The available data are shown in Table IV; they are fully compatible with the predictions. Note also that comparison of Tables I and IV gives $\delta_1 \simeq \delta_2$.

The role which we attribute to charge annihilation in $\bar{K}^-p$ could be tested by noting that this process should be absent in $\bar{K}^-n$, so that $\sigma_{inel}^B(\bar{K}^-n)$ should be mainly hyperon annihilation. The lack of data does not allow us to compare with experiment. As shown by Tabel I, however, the theoretical predictions on total cross sections show that $\sigma_T(\bar{K}^-n) - \sigma_T(K^+n)$ is smaller than $\sigma_T(\bar{K}^-p) - \sigma_T(K^+p)$ by about 2 mb, which is very close to $\delta_1 \simeq \delta_2$.

2) pp and pp Scattering

We follow the same procedure as for Kp scattering, first determining the parameters for G_∞ (assumed to be of Gaussian form). From extrapolated fits to the pp total cross sections [26], we take the value

$$\sigma_T^{(\infty)} = 38.25 \text{ mb.}$$

Values of the logarithmic slope β of pp and $\bar{p}p$ elastic scattering are plotted in Fig. 6, the abscissae being the same as for Kp (Fig. 1). The pp points come from fits of the form $\exp(a + \beta t + \gamma t^2)$ to $d\sigma_{el}/dt$ [24, 27, 28, 29]. Comparison of the forward point deduced by extrapolation with the optical theorem value $(1 + \alpha^2)\sigma_T^2/16\pi$ for reasonable α 's, shows the ref. 27 extrapolation coming out too high and those of refs. 24 and 28 somewhat low. This would imply a tendency for the ref. 27 slopes to be too high and those of refs. 24 and 28 to be too small. The β values for $\bar{p}p$ shown on Fig. 6 come from fits to the $d\sigma_{el}/dt$ data to the form $\exp(a + \beta t)$ [19, 24, 30]. Since the plots of $\ln(d\sigma_{el}/dt)$ curve downwards for increasing $|t|$ this has the effect of yielding too large a value for the slope. Note that the 15.91 GeV/c $\bar{p}p$ point has been corrected by imposing the optical theorem forward point.

The evidence displayed in Fig. 6 would immediately suggest that β_∞ should be taken in the range 10-12 (GeV/c)⁻², probably

below 12 since the $\bar{p}p$ slope appears to be still falling at measured energies. Recent results from Serpukhov [31] indicate that pp slopes are still increasing up to 70 GeV/c. In the region 50-70 GeV/c, β oscillates around 11.4 (7 energies, statistical error $\sim .1$ per energy, systematic error $\sim .3$). An effective one Regge pole fit to all Serpukhov data, covering the region 12-70 GeV/c, gives a trajectory slope $\alpha'_p = 0.40 \pm 0.09$, which corresponds to $\beta(50 \text{ GeV/c}) - \beta(25 \text{ GeV/c}) = 0.54$. Although there would be no difficulty of principle for incorporating an infinitely increasing β in our scheme, we do here as for total cross sections and concentrate our analysis on the region where both pp and $\bar{p}p$ data are available. We fix on the value

$$\beta_\infty = 11.2 (\text{GeV/c})^{-2}.$$

This is the lowest value allowed on the Gaussian overlap model with $\sigma^{(\infty)} = 38.25 \text{ mb}$ and corresponds to

$$P = 1.0, \quad A = 6.33 (\text{GeV/c})^{-2}.$$

Lower values can of course be obtained with other shapes, e.g. the Durand-Lipes form. The resulting $d\sigma^{(\infty)}/dt$ is shown together with a sample of experimental points [24] in Fig. 7. Here again the effect of the real part of the amplitude is seen to be still sizable.

To determine c and ν the $\bar{p}p$ total cross section data [2] were fitted to the form

$$\sigma_T(\bar{p}p) = \sigma_T^{(1)} = \sigma_T^{(\infty)} + 68.58 q_L^{-0.65} = 38.25 + 68.58 q_L^{-0.65}$$

where σ_T is measured in mb and q_L in GeV/c. (The Serpukhov data were not included, but in this case they are well fitted by the above formula). Note that a similar fit to the pp cross section gives

$$\sigma_T(pp) = 38.25 + 16.02 q_L^{-0.99}$$

illustrating that this cross section is much closer to the asymptotic limit than $\sigma_T(\bar{p}p)$.

We now proceed to the determination of δG as for the Kp discussion. From the $\bar{p}p$ data of Foley et al. [19, 24] and our assumed form for $d\sigma_{el}^{(\infty)}/dt$ we compute the $Y(t)$ and fit them to our trial form (43) by varying the parameters H , B and B' . The results are shown in Fig. 8 in which are displayed the experimental points and four representative fits I, II, III and IV. The parameters for these are listed in Table V; also the parameters $J_{A,B}$ from which our predictions for $\sigma_{inel}^{A,B}$ are to be calculated, eq. (46). In Fig. 9, the $\Phi_{A,B}$ which determine the $\delta G_{A,B}$ are plotted as functions of b^2 for parameter choice III, separation methods 1, 2 and 3. For all four choices I to IV the $\Phi_{A,B}$ on methods 2 and 3 are very similar.

Table V shows that for the pp case, B, the range of the B-type processes, is of the same order as A, the range of the asymptotic A-type processes. This is in marked contrast to what we found for the $K^\pm p$ reactions. Comparing Tables III and V, we see that $A(\bar{p}p) \simeq B(\bar{p}p) \simeq 6 \text{ (GeV/c)}^{-2}$ is of the same order as $B(\bar{K}^-p)$, the range of $\bar{K}^-p$ annihilation. It is $A(Kp)$ which is smaller than this common value by about a factor 2; the difference also appears through the fact that β_∞ is $\sim 6 \text{ (GeV/c)}^{-2}$ for Kp and about twice as large for pp. Hence, baryon and hyperon annihilations are similar in range, whereas Pomeron exchange for Kp is quite different from pp. We recall here that for the $\pi^\pm p$ case, Henzi found $A \simeq B \simeq 5 \text{ (GeV/c)}^{-2}$ in our present notation. Comparison of Tables III and V further shows that J_A is systematically larger for $\bar{p}p$ than for $\bar{K}^-p$. This is physically understandable on two grounds; the larger value of $\delta\sigma/\sigma_T^{(\infty)}$ for $\bar{p}p$, and the fact that $G_\infty(pp)$ is closer than $G_\infty(Kp)$ to the unitarity limit as shown by the P values.

Our predictions for σ_{inel}^B may be compared with the measured values of σ_{ann} at 3.3, 5.7 and 7.2 GeV/c [32, 33, 1]. The relevant data is listed in Table VI, the values of $\delta\sigma$ in particular being 37.2, 25.4 and 21.8 mb respectively. The average of our various determinations of the parameter J_B (Table V) is 1.0 ± 0.2 so that we predict

$$\sigma_{\text{ann}} \simeq \sigma_T(\bar{p}p) - \sigma_T^{(\infty)} = \delta\sigma$$

which fits the observations very well. This feature had already been noted by Kokkedee [4]. Since $\delta\sigma$ at 20 GeV/c is still of the order of 10 mb, we predict $\sigma_{\text{ann}}(20 \text{ GeV/c}) \simeq 10 \text{ mb}$, a strikingly large value. The new $\sigma_T(\bar{p}p)$ measurements [17] together with the assumption $\sigma_T^{(\infty)} = 38.2 \text{ mb}$ would imply a corresponding figure of about 5 mb at 50 GeV/c.

There have been indications [34, 35] for the existence of a dip in the $\bar{p}p$ differential cross section at $|t| \simeq 0.6 (\text{GeV/c})^2$. We have not used the wider angle data in our fits. Nonetheless, our solutions do predict diffraction zeros at momentum transfers depending on the beam momentum; for example, at 6 GeV/c, solutions I and III imply zeros at $|t| \simeq 0.65$ and II and IV at $|t| \simeq 0.8$. The zeros move further out at higher beam momenta. Figure 10 shows $d\sigma_{\text{el}}^{(1)}/dt$ corresponding to our fit I to $Y(t)$ compared to the data at 8.9 GeV/c.

To gauge sensitivity to the assumed asymptotic form for $\bar{p}p$ versus pp , we re-computed and fitted $Y(t)$ starting with the Durand-Lipes form for G_∞ , eqs. (15) and (16). The parameters λ and μ were chosen

$$\lambda = 1.43, \quad \mu = 0.914 (\text{GeV/c})^2$$

so as to achieve the same values for $\sigma_T^{(\infty)}$ and β_∞ , 38.25 mb and $11.2 (\text{GeV}/c)^{-1}$, as for the Gaussian overlap discussion above. The resulting change in $d\sigma_{el}^{(\infty)}/dt$ is shown in Fig. 11, the effect on $Y(t)$ being to push it upwards by an amount rising to ~ 0.05 for the larger values of $|t|$ considered. The new $Y(t)$ was fitted to our parametric formula as before with a similar range of results. The parameters for two typical fits are listed in Table V, labelled cases V and VI, together with the results for Gaussian asymptotics. It will be seen that the predictions for non-asymptotic cross sections and the range of B-type processes are essentially unaltered. The feature that H large implies J_B large was true also for the Gaussian G_∞ ; here again method 1 ceases to be reasonable for the larger H value. Fit V implies a diffraction zero for $\bar{p}p$ lying at $|t| \simeq 0.65 (\text{GeV}/c)^2$ for $q_L = 6 \text{ GeV}/c$ and at $|t| \simeq 0.75 (\text{GeV}/c)^2$ for $q_L = 18 \text{ GeV}/c$. Fit VI implies that there is no diffraction zero. The conclusion of our discussion of the modified G_∞ is that the data considered here is rather insensitive to this quantity, once $\sigma_T^{(\infty)}$ and β_∞ are fixed.

IV. Concluding Remarks

We have analysed the $\bar{K}^-p$ versus K^+p and $\bar{p}p$ versus pp differences in σ_T and $d\sigma_{el}/dt$ by attributing them to the shadow effect of baryon, hypercharge and charge annihilation channels. Although model assumptions were needed, we found reasonable results for the total annihilation cross sections and the range of annihilation processes. These results are qualitatively the same for various model assumptions. They are quite sensitive, however, to the input data of our analysis, which are σ_T and $d\sigma_{el}/dt$ for beam momenta $5 \lesssim q_L \lesssim 20$ GeV/c. We found that the data available for the latter quantities, especially for $K^\pm p$ elastic scattering, are far from satisfactory and contain a number of defects which make new experiments in the same energy range highly desirable.

Our analysis has used simple and familiar assumptions concerning the asymptotic behaviour of $d\sigma_{el}/dt$ and σ_T at infinite energy. Recent Serpukhov data in the range $20 \lesssim q_L \lesssim 70$ GeV show that the actual situation may be more complicated than assumed until now. It is therefore important to stress that the approximations in eqs. (32, 33) have only been made because the present data do not warrant a more precise treatment. The essentials of our scheme remain unaffected for more complicated s -dependences of $F^{(0)}$ and $F^{(1)}$, and the calculations can be carried through for any particular s where $G^{(0)}$ and $G^{(1)}$ are available from σ_T and elastic scattering. At least for Kp , good polarization data may also allow a future treatment of the spin, similar to what has been done for πp [6].

Our analysis predicts total annihilation cross sections and ranges. The former are directly measurable for baryon annihilation ($\bar{p}p \rightarrow$ mesons) and hypercharge annihilation (which, except for minor corrections, is $\bar{p}p$ or $\bar{K}^-p \rightarrow$ hyperon or $Y^* +$ non-strange mesons). The concept of charge annihilation which we have introduced is not sufficiently well defined to allow direct experimental tests. This is clearly a shortcoming which requires further work. It seems that the best approach would be to investigate the correlations between

charge and longitudinal velocity for the secondaries of high energy collisions, and to search for traces of the additional charge exchange channels which we have called charge annihilation.

From the theoretical side, as models for inelastic collisions with $n \geq 3$ particles in the final state are being developed, e.g. multi-Regge or Veneziano models, one can apply them to the various types of annihilation processes and calculate their contribution to the overlap function, which should be equal to the function $\delta G_B(b,s)$ determined by our analysis. In particular, the models will lead to an s-dependence of δG_B which will result from an interplay of two effects, firstly the average value of the Regge trajectories $\alpha(t)$ dominating annihilation, and secondly the increase with s of the final state multiplicities. Baryon annihilation would be the simplest case on which to try this theoretical study.

One can also ask whether the annihilation range predicted by our analysis could be tested directly against experiment. This requires the experimental determination of the range of the overlap function δG_B for annihilation processes. Mainly from the work of Michejda on the complete overlap function of πp collisions [36] we know that this is very difficult, because the final state wave functions must contain either rapidly varying phase factors or subtle correlations between particles, both of which escape experimental observation so far.

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APPENDIX

K-Matrix and Eikonal Models for Combined Shadow of Absorptive Channels

Our object is to construct plausible models which are explicitly unitary and exhibit the mutual influence of competing inelastic channels. We picture an underlying many-channel situation which leads to purely absorptive scattering in channel 1, the elastic channel of interest. We further suppose the inelastic channels to be divisible into two disjoint and exhaustive groups. In the first model, we explicitly treat an n -channel situation and assume that the $n \times n$ K-matrix is additive. In the second model, we assume that the elastic S-matrix is multiplicative, as would be the case for scattering through a block of material containing two kinds of absorbers.

We normalise so that the S-, T- and K-matrices are connected via the equations

$$S = I + 2iT = (I + iK)(I - iK)^{-1} \quad (\text{A.1})$$

where I is the unit matrix. We shall write our equations in a notation appropriate for two-body channels of spinless particles with fixed angular momentum, but this is not an essential restriction. By unitarity, the total overlap function G is defined by

$$G = \sum_{n \neq 1} |S_{1n}|^2 = \sum_{n \neq 1} G_n \quad (\text{A.2})$$

Here, we have introduced the subdivision of G into partial overlap functions G_n , the contributions from particular final channels.

K-Matrix Model

For this, we assume the simplest possible form for the K-matrix

$$K = \begin{pmatrix} 0 & K_{12} & K_{13} \\ K_{12} & 0 & 0 \\ K_{13} & 0 & 0 \end{pmatrix} \quad (\text{A.3})$$

We set down the equations as if there were just three channels but the result extends without alteration to the case of n channels where 2 and 3 denote blocks of channels with non-vanishing matrix elements connecting to channel 1. Solving for the S-matrix elements by (A.1), one finds

$$\left. \begin{aligned} S_{11} &= (1 - K_{12}^2 - K_{13}^2) (1 + K_{12}^2 + K_{13}^2)^{-1} \\ S_{12} &= 2i K_{12} (1 + K_{12}^2 + K_{13}^2)^{-1} \end{aligned} \right\} \quad (\text{A.4})$$

and similarly for S_{13} , with of course $|S_{11}|^2 + |S_{12}|^2 + |S_{13}|^2 = 1$. The partial overlap function G_2 is given by

$$G_2 = |S_{12}|^2 = 4 K_{12}^2 (1 + K_{12}^2 + K_{13}^2)^{-2} \quad (\text{A.5})$$

One can also introduce the hypothetical quantity g_2 for the overlap if channel 3 were absent and further express K_{12} in terms of this:

$$g_2 = 4 K_{12}^2 (1 + K_{12}^2)^{-2} \quad (\text{A.6})$$

$$K_{12}^2 = -1 + (2/g_2) [1 - (1 - g_2)^{1/2}] \quad (\text{A.7})$$

We similarly introduce G_3, g_3 .

The model proceeds by making assumptions on the form of g_2 and g_3 . In our application, we identify channel 2 with the persisting A-type processes and channel 3 with the evanescent B-type processes. Our physical assumption is therefore

$$g_2 = G_\infty \quad (\text{A.8})$$

$$g_3 = \delta g \quad (\text{A.9})$$

One can now set down the resulting formulae in terms of G_∞ and δg for the physical quantities $G_2 = G_A$ and $G_3 = G_B$. For simplicity, we retain only linear terms in δg . Thus,

$$G_A = G_\infty - \frac{1}{4} [1 - (1 - G_\infty)^{1/2}]^{-1} G_\infty^2 \cdot \delta g \quad (\text{A.10})$$

$$G_B = \frac{1}{4} [1 - (1 - G_\infty)^{1/2}]^{-2} G_\infty^2 \cdot \delta g \quad (\text{A.11})$$

giving

$$\begin{aligned} G &= G_A + G_B = G_\infty + \delta G = \\ &= G_\infty + \frac{1}{4} [1 - (1 - G_\infty)^{1/2}]^{-2} (1 - G_\infty)^{1/2} \cdot G_\infty^2 \cdot \delta g \end{aligned} \quad (\text{A.12})$$

This leads immediately to model (2) for the partition of δG

$$\delta G_A = G_\infty - G_A = [1 - (1 - G_\infty)^{1/2}] \cdot (1 - G_\infty)^{-\frac{1}{2}} \cdot \delta G \quad (\text{A.13})$$

$$\delta G_B = G_B = (1 - G_\infty)^{-1/2} \cdot \delta G \quad (\text{A.14})$$

Eikonal Model

Here we invoke the notion of two co-existing absorbing mechanisms in a slab. Then, if n_1 is the intensity of the transmitted beam at thickness x , n_2 and n_3 the number of particles absorbed via the two mechanisms and λ_2 and λ_3 the respective absorption coefficients, we obtain

$$\left. \begin{aligned} n_1 &= \exp \left[-2 (\lambda_2 + \lambda_3) x \right] \\ n_2 &= \lambda_2 (\lambda_2 + \lambda_3)^{-1} \cdot (1 - n_1) \\ n_3 &= \lambda_3 (\lambda_2 + \lambda_3)^{-1} \cdot (1 - n_1) \end{aligned} \right\} \quad (\text{A.15})$$

This suggests the model (we use the same notation as for the K-matrix model discussed above)

$$1 - G = S_{11}^2 = (1 - g_2)(1 - g_3) \quad (\text{A.16})$$

$$G_2 = [1 - (1 - g_2)(1 - g_3)] \ln(1 - g_2) / \ln[(1 - g_2)(1 - g_3)] \quad (\text{A.17})$$

$$G_3 = [1 - (1 - g_2)(1 - g_3)] \ln(1 - g_3) / \ln[(1 - g_2)(1 - g_3)] \quad (\text{A.18})$$

In particular, making the identifications as in equations (A.8, A.9) and assuming δg small, we obtain

$$G_A = G_\infty + \left\{ (1 - G_\infty) + [G_\infty / \ln(1 - G_\infty)] \right\} \cdot \delta g \quad (\text{A.19})$$

$$G_B = - [G_\infty / \ln(1 - G_\infty)] \cdot \delta g \quad (\text{A.20})$$

$$G = G_{\infty} + (1 - G_{\infty}) \cdot \delta g = G_{\infty} + \delta G \quad (\text{A.21})$$

This leads to model (3) for the partition of δG

$$\delta G_A = - \left\{ 1 + G_{\infty} [(1 - G_{\infty}) \ln(1 - G_{\infty})]^{-1} \right\} \cdot \delta G \quad (\text{A.22})$$

$$\delta G_B = - [(1 - G_{\infty}) \ln(1 - G_{\infty})]^{-1} \cdot G_{\infty} \cdot \delta G \quad (\text{A.23})$$

Finally, taking the two models together, the following points may be noted:

- (i) The two models give essentially the same results except where G_{∞} is very close to unity. This can be checked by explicit calculation or by making power series expansions in G_{∞} or $\ln(1 - G_{\infty})$.
- (ii) The two models furnish via equations (A.11, A.12) and (A.20, A.21) prescriptions of how the shape of δG is modified from that of δg by the predominant A-type absorption. If one assumes a simple form for δg they yield explicit parametrizations for the physical δG .

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TABLE CAPTIONS

- TABLE I : Summary of total cross sections for various channels in the KN system.
Figures are from the fits of Dass, Michael and Phillips [22], their solution III. Experimental results at 16 GeV/c from Galbraith et al. [2] are shown in brackets. All cross sections in millibarns, q_L in (GeV/c).
- TABLE II : Values assumed for $\alpha = (\text{Re}F / \text{Im}F)_{t=0}$ for K^+p elastic scattering and the corresponding values of $(d\sigma_{el} / dt)_{t=0}$, q_L in (GeV/c), $d\sigma_{el} / dt$ in $\text{mb}/(\text{GeV}/c)^2$. The result of our extrapolation to $q_L = \infty$ is also shown.
- TABLE III : $\bar{K}^-p$ and K^+p systems. Representative sets of values for the parameters H , B and B' in our fit (43) to the normalized non-asymptotic scattering amplitude $Y(t)$, eq. (42). P and A are the parameters of the asymptotic overlap function, A , B and B' are measured in $(\text{GeV}/c)^{-2}$. The values of $J_B - J_A$ are independent of the method.
- TABLE IV : $\bar{K}^-p$ and K^+p systems. Comparison of our prediction for σ_{inel}^B , the total cross section for the B-type processes, with experimental data on hyperon production in $\bar{K}^-p$ at 6 GeV/c [25] and 10 GeV/c [3], and the difference $\sigma_T(\pi^-p) - \sigma_T(\pi^+p)$ [2]. The first line is from Table I. All cross sections are in millibarns, q_L in GeV/c.

TABLE V : $\bar{p}p$ and pp systems. Representative sets of values I, II, III and IV for the parameters H, B and B' as in Table III but for $\bar{p}p$; also fits V and VI for the Durand-Lipes asymptotic form (15).

TABLE VI : Summary of total cross sections for various channels in the $\bar{p}p$ system, illustrating the importance of baryon and hypercharge annihilation. The latter is partly contained in $\sigma(\bar{p}p \rightarrow \text{kaons} + \text{pions})$ and partly in $\sigma(\bar{p}p \rightarrow \text{hyperons} + \dots)$. Charge annihilation was not estimated but should be small. At 7 GeV/c, the value for $\sigma(\bar{p}p \rightarrow \bar{N}N + \text{pions})$ does not include the contribution from $\bar{p}p \rightarrow \bar{n}n + 2$ or more pions, which was included in the number for $\sigma(\bar{p}p \rightarrow \text{pions})$. The last number for 7 GeV/c is $\sigma(\bar{p}p \rightarrow \text{strange particles} + \dots)$. The data are from Ferbel et al. at 3.3 GeV/c [32], Böckmann et al. at 5.7 GeV/c [33] and T. Ferbel et al. at 7.0 GeV/c [1].

T A B L E I

q_L	6	8	10	12	16	16	20
$\sigma_T(\bar{K}^-p)$	24.2	23.1	22.4	21.9	21.1	(21.3 $\pm$ 0.4)	20.6
$\sigma_T(\bar{K}^-n)$	22.0	21.2	20.7	20.3	19.7	(20.3 $\pm$ 0.5)	19.4
$\sigma_T(K^+p)$	17.1	17.2	17.3	17.3	17.3	(17.0 $\pm$ 0.1)	17.2
$\sigma_T(K^+n)$	17.7	17.7	17.8	17.8	17.7	(17.4 $\pm$ 0.4)	17.7

T A B L E II

	$K^+ p$				$K^\pm p$	$\overline{K}^- p$				
q_L	6.8	9.8	12.8	14.8	∞	7.2	9.8	10.12	11.88	15.31
$\left(\frac{d\sigma_{el}}{dt}\right)_{t=0}$	16.6 ± 0.3	16.0 ± 0.3	16.0 ± 0.3	16.0 ± 0.3	15.2	28.4 ± 0.3	26.3 ± 0.3	25.8 ± 0.3	24.7 ± 0.3	23.4 ± 0.3
$(\alpha)_{t=0}$	0.32 ± 0.05	0.23 ± 0.05	0.23 ± 0.05	0.23 ± 0.05	0	≈ 0	≈ 0	≈ 0	≈ 0	≈ 0

T A B L E III

PARAMETER CHOICE	H	B	B'	J_B^1	J_B^2	J_B^3	$J_B - J_A$
I P = 0.90	1.24	6.0	3.32	1.82	1.00	1.00	0.91
II A = 3.32	0.6	6.7	2.4	1.32	0.99	1.00	0.91
III	0.33	7.3	2.0	1.11	1.00	1.01	0.89
IV P = 0.85	0.1	8.0	1.5	0.94	1.01	1.02	0.87
A = 3.53							

Average $J_B \approx 1.1 \pm 0.25$

T A B L E IV

q_L	6	8	10	12	16	20
$\sigma_{\pi}(\bar{K}^-p) - \sigma_{\pi}(K^+p)$	7.1	5.9	5.1	4.6	3.8	3.3
$\sigma(\bar{K}^-p \rightarrow \Lambda + \text{pions})$	3		2.3 ± 0.2			
$\sigma_{\pi}(\pi^-p) - \sigma_{\pi}(\pi^+p) = \delta_2$	2.3 ± 0.4	2.4 ± 0.4	1.7 ± 0.4	1.7 ± 0.4	1.7 ± 0.4	1.1 ± 0.4
σ_{inel}^B	7.8 ± 1.8	6.5 ± 1.5	6.3 ± 1.3	5.1 ± 1.2	4.2 ± 1.0	3.6 ± 0.8

T A B L E VI

q_L	3.3	5.7	7.0
$\sigma_T(\bar{p}p)$	75.4 ± 2.0	63.6 ± 1.4	60.0 ± 2.0
$\delta\sigma = \sigma_T(\bar{p}p) - \sigma_\infty$	37.2	25.4	21.8
$\sigma(\bar{p}p \rightarrow \bar{N}N + \text{pions})$	19.2 ± 2.0	24.8 ± 2.0	19.5 ± 2.8
$\sigma(\bar{p}p \rightarrow \text{pions})$	30.9 ± 3.0	22.5 ± 2.0	23.6 ± 3.4
$\sigma(\bar{p}p \rightarrow \text{kaons} + \dots)$	2.9 ± 0.2		
$\sigma(\bar{p}p \rightarrow \text{hyperons} + \dots)$	0.44 ± 0.05		2.7 ± 0.7

FIGURE CAPTIONS

Figure 1 : Logarithmic slopes β from fits to $d\sigma_{el}/dt$ for K^+p and $\bar{K}^-p$, and extrapolation to infinite energy. K^+ slopes $\beta^{(0)}$ (left) based on data from [19] ; $\bar{K}^-$ slopes $\beta^{(1)}$ (right) from [19, 23, 24] . Fits with forward optical point imposed, full line error bars; without forward constraint, dotted error bars. The two alternatives for $\beta^{(\infty)}$ treated in text are indicated on the central ordinate. Values of P corresponding to various $\beta^{(\infty)}$ for $\sigma_T^{(\infty)} = 17.2$ mb are shown on the side. The dotted lines are drawn by eye.

Figure 2 : Alternative versions of Kp asymptotic cross section $d\sigma_{el}^{(\infty)}/dt$ compared to high energy K^+p data. Asymptotic parameters $P = 0.90$, $A = 3.32$ and $P = 0.85$, $A = 3.53$ (GeV/c) $^{-2}$. Data points are: $\bar{\Phi}$ 14.8, $\bar{I}$ 12.8, $\bar{\Phi}$ 9.8, $\bar{I}$ 6.8 and $\bar{A}$ 7.3 GeV/c [19] . The quoted overall scale uncertainties [19] have been included in the enlarged errors. At $t = 0$, the bar I indicates the range of assumed forward points (Table II).

Figure 3 : Phenomenological values of $Y(t)$ for $\bar{K}^-p$ and fits by formula (43).
Fig. 3a for $P = 0.90$, fits I (---), II (—) and III (.....).
Fig. 3b for $P = 0.85$, fit IV (—).
 Data points are: $\bar{\Phi}$ 15.91 and $\bar{\Psi}$ 11.88 GeV/c [24] , $\bar{I}$ 10.12 GeV/c [23] , $\bar{\Phi}$ 9.0 and $\bar{X}$ 7.2 GeV/c [19] .

Figure 4 : Comparison of $d\sigma_{el}^{(\infty)}/dt$ computed for $P = 0.90$ (full line) and $P = 0.85$ (dashed line) with $\bar{K}^-p$ data; legend as for Fig. 3. The arrow shows the range of forward points assumed (Table II).

Figure 5 : Illustration of alternative possibilities for the functions $\Phi_{A,B}(b^2)$ related to $\delta G_{A,B}$ by eq. (45), for $\bar{K}^-p$.

Fig. 5a, comparison of separation prescriptions 1, 2 and 3 for the δG of fit II.

Fig. 5b, forms given by method 2 for fits I to IV.

Figure 6 : Logarithmic slopes β for pp and $\bar{p}p$. Left, $\beta^{(0)}$ for pp: $\bar{\nabla}$ from [28], $\bar{\Phi}$ from [24], $\bar{\Phi}$ from [27], $\bar{\nabla}$ from [29]. Right, $\beta^{(1)}$ for $\bar{p}p$: $\bar{\Phi}$ from [19], $\bar{\Phi}$ from [24], $\bar{\nabla}$ from [30]; values taken directly from references except for higher $\bar{p}p$ point at 15.91 GeV/c which is obtained imposing the forward optical point. Dotted lines are drawn by eye.

Figure 7 : Asymptotic $d\sigma_{el}^{(\infty)}/dt$ for pp and $\bar{p}p$ compared to a typical set of pp data points for $d\sigma_{el}/dt$: $\bar{\nabla}$ 10.94, $\bar{\Phi}$ 14.93, $\bar{\Phi}$ 19.82, $\bar{\nabla}$ 21.88 and $\bar{\nabla}$ 24.63 GeV/c, all from [24]. The quoted overall scale uncertainty has been included in the enlarged errors. Asymptotic parameters $P = 1.0$ and $A = 6.33 (\text{GeV}/c)^{-2}$.

Figure 8 : Phenomenological values of $Y(t)$ for $\bar{p}p$, calculated from the $d\sigma_{el}^{(\infty)}/dt$ of Fig. 7 and the data points: $\bar{\Phi}$ 12.0, $\bar{\Phi}$ 10.0, $\bar{\nabla}$ 8.9 and $\bar{\nabla}$ 7.2 GeV/c from [19], and $\bar{\nabla}$ 15.91 and $\bar{\nabla}$ 11.8 GeV/c from [24]. The curves depict the fits by formula (43), fits I to IV of Table V: (I) ---, (II) -.-., (III) —, (IV)

Figure 9 : Contributions $\Phi_{A,B}(b^2)$ to non-asymptotic overlap function for $\bar{p}p$ (cf. Fig. 5) corresponding to fit III of Table V. Curves correspond to prescriptions 1, 2 and 3.

Figure 10 : $d\sigma_{el}/dt$ for $\bar{p}p$ corresponding to fit I compared to the data at 8.9 GeV/c from [19] .

Figure 11 : Comparison of $[d\sigma_{el}^{(\infty)}/dt]^{1/2}$ for the Durand-Lipes (15) and Gaussian form (12) for G_{∞} (same $\sigma_T^{(\infty)}$ and $\beta^{(\infty)}$).

Figure 1

Figure 2

Figure 3a

Figure 3b

Figure 4

Figure 5a

Figure 5b

Figure 6

Figure 7

Figure 8

Figure 9

Figure 10

Figure 11