AN APPLICATION OF THE ABSORPTION MODEL TO THE REACTIONS
\[ p\bar{p} \rightarrow Y\bar{Y} \]

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ABSTRACT

We present a set of calculations of the reaction \[ p\bar{p} \rightarrow Y\bar{Y} \] in the antiproton lab. momentum range from 3 GeV/c to 7 GeV/c. The process is assumed to take place through \( k^* \) exchange only, and the calculations are performed within the framework of the absorption model (the Distorted Wave Born Approximation). We obtain a good description of the angular dependence of the cross-section for the three channels \( \Lambda\bar{\Lambda}, \Lambda^0\bar{\Lambda}^0 \) and \( \Sigma^+\bar{\Sigma}^- \) in the whole energy range. However, the energy dependence is wrong, and appears the worst for the channel \( \Lambda\bar{\Lambda} \), where the angular dependence is best described.

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1. INTRODUCTION

The process antiproton-proton producing a pair antihyperon-
hyperon has recently been investigated in experimental \textsuperscript{1},\textsuperscript{2} and
theoretical \textsuperscript{3}-\textsuperscript{7} papers.

Experiments are characterized by a strong forward peaking of
the antihyperon relative to the direction of the incoming antiproton,
indicating predominance of higher partial waves characteristic for what
one calls peripheral processes.

Now it is well known that by calculating a Born diagram for
the process with the exchange of a $K$ or a $K^*$ meson, one gets a
contribution from low partial waves exceeding the bounds set by unitarity.
This results in an angular distribution for the produced particle which
shows a too weak angular dependence. As, on the other hand, Born diagrams
are easily calculable, and give a simple intuitive picture of what is
happening, it was then natural to seek methods by which one would unita-
rise the Born approximation \textsuperscript{8},\textsuperscript{9}, or at least cut down the contribution
of low partial waves.

The method we shall follow here bears on this last point.
Originating from Sopkovitch \textsuperscript{4}, it was later elaborated \textsuperscript{5},\textsuperscript{6},\textsuperscript{10} and
is usually called the absorption model. The damping of the low partial
waves is obtained by including the elastic interaction in the initial
and final states, which shows up as an "absorption" factor modifying
each partial wave of the helicity amplitudes. This absorption factor
is directly related to the elastic scattering data. For theoretical
discussions of that procedure we refer to Refs. \textsuperscript{9}-\textsuperscript{12}.

Earlier theoretical works, either based on the Born diagrams \textsuperscript{3}
or on the Sopkovitch type of calculations \textsuperscript{4}-\textsuperscript{7} indicate that in order to
fit the angular distribution, it is necessary to assume that $K^*$ meson
exchange is predominant. We therefore assume the process to take place
as in Fig. 1.
Our calculations differ from the previous ones by using, together with the usual vector coupling of $K^*$ to the baryon vertex, a tensor coupling (magnetic moment type), and by the systematical investigation of the behaviour of the cross-sections with the variation of the incoming antiproton energy.

In Section 2 we write down the most general form of the interaction necessary to obtain the diagram of Fig. 1, show how we make the partial wave expansion of the helicity amplitudes and how we correct them for the interaction in the initial and final states. For clarity the expressions for the amplitudes together with their partial wave expansion are put in an Appendix.

In Section 3 we present and discuss the numerical results of the calculations for $p\bar{p} \rightarrow \Lambda\bar{\Lambda}$, $p\bar{p} \rightarrow \Sigma^+\bar{\Sigma}^+$ and $p\bar{p} \rightarrow \Lambda\bar{\Sigma}^0 + \Lambda\bar{\Sigma}^0$, with antiproton lab. momentum of 3 GeV/c, 3.6 GeV/c *), 5.7 GeV/c and 7 GeV/c (this last energy only for the $\Lambda\bar{\Lambda}$ channel).

The results illustrate once more the failure of the model to describe correctly the energy dependence of the cross-section when the exchanged particle has a spin different from zero 12). On the other hand, the angular dependence is very well described for $p\bar{p} \rightarrow \Lambda\bar{\Lambda}$, rather well for $p\bar{p} \rightarrow \Lambda\bar{\Sigma}^0 + \Lambda\bar{\Sigma}^0$, and is in reasonably good agreement with experiment for $p\bar{p} \rightarrow \Sigma^+\bar{\Sigma}^+$.

*) In order to have better statistics, the angular distribution at 3.6 GeV/c includes also events from the same reaction at 4 GeV/c.
2. **CALCULATIONS**

To describe the process in Fig. 1 we start from the following effective Lagrangian which has the most general form

\[ \mathcal{L}(x) = \left[ i q_V \bar{\psi}_N(x) \gamma_\mu \psi_N(x) + \frac{g_{\pi}}{m_{\pi}^0 m_N} \partial_\mu \left( \bar{\psi}_N(x) \sigma_{\mu \nu} \psi_N(x) \right) \right. \]

\[ + i \frac{g_S}{m_N^2 m_Y} \partial_\mu \left( \bar{\psi}_N(x) \psi_N(x) \right) \Gamma_\mu(x) + \text{Herm. conj.} \]

(2.1)

Here \( \psi_N(x) \) represents the hyperon field, \( \psi_N(x) \) the nucleon field, \( \Gamma_\mu(x) \) the \( K^* \) field (creating \( K^{*+} \) and destructing \( K^{*-} \)), \( g_V, g_T \) and \( g_S \) are the vector, tensor and scalar coupling constants of \( K^* \) to the nucleon-hyperon vertex and \( m_N \) and \( m_Y \) are the masses of the nucleon and hyperon.

We shall further denote by \( a^\mu, b^\mu, c^\mu \) and \( d^\mu \) the four-momenta for \( \bar{\nu}, p, \bar{\Sigma} \) and \( \Sigma \), respectively as indicated in Fig. 1, and we define

\[ e^\mu = b^\mu - d^\mu = c^\mu - a^\mu \quad , \quad t = - e^\mu e_\mu. \]

The calculations will be made in the centre-of-mass system and we shall assume the same mass for the hyperon and antihyperon. (When we calculate the process \( p \bar{p} \to \Lambda \Sigma^0 + \bar{\Lambda} \Sigma^0 \) we will then use as the hyperon mass the mean of the masses of \( \Lambda \) and \( \Sigma^0 \) but we believe that the error so induced is unimportant.) Then in the centre-of-mass system the four particles have the same energy \( E \), and we denote by \( s \) the total energy of the incoming particles squared : \( s = 4E^2. \)
Still in the centre-of-mass system the magnitude of the three momentum is denoted by $q$ for the proton and antiproton, by $q'$ for the hyperon and antihyperon while $\theta$ is the angle between the proton and hyperon momenta.

$$\sigma_{\mu} \quad \text{and} \quad \sigma_{\mu\nu} = \frac{1}{2i} (\gamma_{\mu} \gamma_{\nu} - \gamma_{\nu} \gamma_{\mu})$$ are the ordinary Dirac operators acting on the four dimensional spinor space when we normalize the free-particle spinors to $\bar{u}u = 2m$ where $m_x$ is the mass of the baryon. The production amplitudes $M$ for the process $p\bar{p} \to YY$ will be normalized so that the differential cross-section is given by

$$\frac{d\sigma}{dt} = \frac{4}{64s q^2} \frac{1}{4} \sum |M|^2$$

where the summation is over the 16 helicity amplitudes of the process. The mass of $K^*$ will be denoted by $m_x$.

With these notations the Lagrangian (2.1) gives the following expression for the amplitudes

$$M = \bar{\psi}_x (\mathbf{q}) \left( \gamma_v \gamma^\nu + \frac{\mathbf{q} \cdot \mathbf{e}}{\mathbf{q}^2} \left[ \frac{\sigma \mathbf{e}_\nu}{m_N + \mathbf{m}_y} \mathbf{e}_\nu + \frac{\mathbf{m}_\nu + \mathbf{m}_y}{m_N + \mathbf{m}_y} \mathbf{e}_\nu \right] \right) \psi_x (\mathbf{k})$$

where $\mathbf{v}$ represents the negative energy Dirac spinors. Using the properties of the charge conjugation operator $C$ (the superscript $t$ denotes transposed):

$$C^{-1} \bar{\psi}_x (-k) = \bar{\mu}_x (k)$$

$$\bar{\psi}_x (-k) = - \mu_x^t (k) C^{-1}$$

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and \[ C^{-1} \gamma_r C = - \gamma^r \]

the expression (2.2) is easily transformed into

\[
\mathbf{M} = - \bar{\mu}(c) \left\{ g_V \gamma_5 - \frac{g_T}{m^2_{w^+w^-}} \sigma^{\mu\nu} g_5 e_s - \frac{i g_S}{m^2_{w^+w^-}} e_s \right\} \mu(a) \\
\frac{\delta_{\mu\nu} + \frac{e_\mu e_\nu}{m^2}}{m^2-t} \bar{\mu}(d) \left\{ g_V \gamma_5 + \frac{g_T}{m^2_{w^+w^-}} \sigma^{\mu\nu} e_s + \frac{i g_S}{m^2_{w^+w^-}} e_s \right\} \mu(b)
\]

(2.3)

We are going to make the assumption that the baryon always can be treated as being on the mass shell, and by use of the Dirac Eq. (2.3) can be written as

\[
\mathbf{M} = G_{VV} \bar{\mu}(d) \gamma_5 \mu(b) \bar{\mu}(c) \gamma_5 \mu(a) \frac{1}{m^2-t} \\
+ G_{4S} \left[ \bar{\tau}(d) \gamma_5 \mu(b) \bar{\mu}(c) \gamma_5 \mu(a) + \bar{\mu}(d) \mu(b) \bar{\mu}(c) \gamma_5 \mu(a) \right] \frac{1}{m^2-t} \\
+ G_{SS} \bar{\mu}(d) \mu(b) \bar{\mu}(c) \mu(a) \frac{1}{m^2-t}
\]

(2.4)

the quantities \( G_{VV}, G_{4S} \) and \( G_{SS} \) are defined by:

\[
G_{VV} = -(g_V + g_T)^2 \\
G_{4S} = \frac{4Eg_4}{m^2_{w^+w^-}} \\
G_{SS} = -2(g_T + g_4)g_T^2 - \frac{g_T^2}{m^2_{w^+w^-}} \left( \frac{m^2_{e_b-e_a}}{m^2_{w^+w^-}} \right)^2 \\
\left( \frac{g_T^2}{m^2_{w^+w^-}} \right)^2 - \frac{1}{m^2} \left[ (m^2_{w^+w^-}) g_T + g_4 \frac{1}{m^2_{w^+w^-}} \right]^2
\]

(2.5)
One can see that for a reasonable scalar coupling the terms containing \( g_S \) are negligible. We shall, therefore, omit them from now on, setting \( g_S = 0 \).

The calculations of the helicity amplitudes are now straightforward, but lead to rather long expressions which we give in an Appendix together with their expansion in terms of the rotation functions \( \tilde{d}_{\lambda m}^j \).

Each helicity amplitude is expressed in the form

\[
\langle \lambda_c \lambda_d | M | \lambda_a \lambda_b \rangle = \sum_{j^m, m, \lambda} \bar{B}^j(s_z z) \bar{C}^j_{\lambda \lambda} (x) + \sum_{j^m, \lambda} \frac{(2j+1)}{2} C^j_{\lambda \lambda} (x) C^j_{\lambda \lambda'} (z)
\]

(2.6)

where

\[
\begin{align*} 
\chi &= \cos \theta \\
\xi &= 1 + \frac{m^2 + (q - q')^2}{2qq'} \\
\lambda &= \lambda_c - \lambda_b \\
\lambda' &= \lambda_c - \lambda_d
\end{align*}
\]

and

\[
j^m_{\lambda m} = \max(|\lambda|, |\lambda'|)
\]

The \( c^j_{\lambda \lambda'} \) are rotation functions of the second kind \(^{12},^{13}\) defined through the relation

\[
\left( \frac{1 - \lambda}{2} \right) \left( \frac{1 + \lambda}{2} \right) \frac{1}{2 - \chi} = \sum_{j^m, \lambda} (2j+1) C^j_{\lambda \lambda'} (z) C^j_{\lambda \lambda} (x)
\]

(2.7)

Their properties and recurrence formulae for their calculations are discussed elsewhere \(^{12}\).
Each term in these expansions is now modified by a factor $k(j)$ which is determined by the elastic scattering in the initial and final states. If one assumes the elastic scattering cross-section to be described by

$$\frac{d\sigma_{el}}{dt} = \frac{\sigma_T^2}{16\pi} \frac{A_t}{t}$$

(2.8)

where $t$ is the momentum transfer squared, one is led to the form

$$R(j) = \left[ \left( 1 - C_i \epsilon_i \gamma_i j \right) / \left( 1 - C_f \epsilon_i \gamma_i j^2 \right) \right]^{1/2}$$

(2.9)

Here

$$C_i = \frac{\sigma_T^i}{4\pi A_i} \quad \gamma_i = \frac{l}{2q_2 A_i} \quad \gamma_f = \frac{l}{2q_2 A_f}$$

(2.10)

$A_i (A_f)$ represent the width of the diffraction peak for the elastic scattering in the initial (final) state, as represented in (2.8) and $\sigma_T^i (\sigma_T^f)$ represent the corresponding total cross-section.

From experimental data on $p\bar{p}$ elastic scattering, we find $C_i = 1$, and the following values for $A_i$

- antiproton lab. momentum $P_{lab.} = 3$ GeV/c $A_i = 16$
- $P_{lab.} = 3.6$ GeV/c $A_i = 15$
- $P_{lab.} = 5.7$ GeV/c $A_i = 13$
- $P_{lab.} = 7$ GeV/c $A_i = 12$
We know, of course, nothing about the elastic scattering of $\bar{Y}Y$, and we chose in a first set of calculations the same values as in the initial state. A discussion about this assumption will be made in the next Section.

Let us note that the values $C_i = C_f = 1$, if they do not ensure necessarily unitarity, at least prevent a violation of the boundedness of partial waves dictated by unitarity when $s$ goes to infinity $^{12})$.

3. **COMPARISON WITH EXPERIMENT**

We performed numerical calculations of the processes $p\bar{p} \rightarrow \Lambda\bar{\Lambda}$, $p\bar{p} \rightarrow \Sigma^+\bar{\Sigma}^-$ and $p\bar{p} \rightarrow \Lambda\bar{\Sigma}^0 + \bar{\Lambda}\Sigma^0$ within the range 3 to 7 GeV/c for the incoming antiproton lab. momentum. It should be remembered that in calculating the amplitudes we assumed equal masses for the two produced particles, for the $\Lambda\bar{\Sigma}^0 + \bar{\Lambda}\Sigma^0$ cross-section we therefore made the approximation $m_{\chi\sigma} = m_{\Lambda} = 1.15$ GeV.

The coupling constants of the $K^*$ to the nucleon-hyperon vertex are unknown, and at each energy we used them as free parameters to fit the angular distribution.

We first tried different values for the ratio $g_T/g_V = R$. If one believes in SU$_3$, the ratio $g_T/g_V$ should be the same for the $K^*\Lambda\bar{\Lambda}$ vertex as for the $\Omega NN$ vertex. This gives from the $\Omega$ photon analogy for the $\Omega NN$ vertex $^{15}$) $R = 3.7$. We found that for values of $R$ in the range from 0 to 1 we could fit the angular distribution of $\Lambda\bar{\Lambda}$ events extremely well, whereas for $R = -1$ or $R = 2$ the forward peak became too broad. Thus within the $K^*$ exchange picture SU$_3$ symmetry appears seriously broken unless it is the $\Omega$ photon analogy or the absorption model which cannot be trusted.
On the other hand, the results are rather insensitive to variation of \( R \) in the range 0 to 1 and we fixed \( g_\pi = g_\pi \) over the whole energy range, leaving one multiplicative coupling constant for each process to fit at each energy \(^*)\). We shall call it the "effective coupling constant". If the model were able to describe well the energy variation of the cross-section the effective coupling constants would be the same at all energies. Thus their variation is a measure of the inability of the model to account for the energy dependence of the cross-sections.

A further problem was to choose among the experimental data. As can be seen from Fig. 2, the data taken from Ref. \(^1\) and Ref. \(^2\) are sometimes in conflict. We usually compare our calculations with the data of Ref. \(^1\) except for \( P = 7 \) GeV/c where only the Yale data are available. The error limits that we put on the effective coupling constants are not to be taken too seriously as they are determined only from the errors in the experimental data used. We recall also the uncertainty resulting from the choice of the \( YY \) elastic parameters (we took them equal to the \( pp \) elastic parameters which also has rather big uncertainties) and to our ignorance of the ratio \( g_T/g_\pi \).

Figures 3 and 4 show the results for \( pp \to \bar{\Lambda} \Lambda \). One sees that the angular distribution is extremely well fitted and this through the whole energy interval. The coupling constants we used to obtain this fit are plotted in Fig. 7, from which it is evident that the model is quite inadequate to describe the rapid fall of the experimental cross-section with energy.

\(^*)\) Using \( g_T = 0 \) and \( (2\Lambda_\pi q^2)^{-1} = (2\Lambda_\pi q^2)^{-1} = 0.26 \) as in Ref. \(^7\) we find at 3.6 GeV/c \( \frac{g^2 K \bar{N} / 4\pi}{4\pi} = 2.8 \pm 0.3 \), in reasonable agreement with the value \( 3.05^{+0.9}_{-0.6} \) of Ref. \(^7\).

The lower value 2.2 obtained in Ref. \(^6, b\) is due partly to the use of a somewhat smaller value for the \( A \)'s and partly to the asymptotic approximation. The \( K \mid \lambda - \lambda' \mid \) Bessel functions appear to be a bad approximation of the \( e_r^\lambda \Lambda \Lambda' \) functions of Ref. \(^6, b\) because \( z \) is relatively far from 1. Also the "extraordinary" terms \( \Theta \) the ones involving \( B^\lambda (s,2) \) in (2.6) are disregarded in Ref. \(^6, b\).
We now turn to the process $p\bar{p} \rightarrow \Sigma^+ \bar{\Sigma}^+$. As can be seen from Fig. 5, the angular distribution for this reaction is rather different from the distribution of the $\Lambda\bar{\Lambda}$ events. Apart from the peak at very low momentum transfer, which is at least as steep as the peak for the $\Lambda\bar{\Lambda}$ case, there appears a number of events at higher $|t|$ values. We call those "weak peripheral events" as they are not of statistical nature, appearing only for $|t|$ values lower than $|t_{\text{max}}|/2$.

In the framework of the model, it appears very difficult to explain these weak peripheral events: increasing the tensor coupling will destroy the narrow peak. The additional exchange of a heavy particle with isospin $\frac{2}{3}$ cannot be considered because this weak peripherality also shows up in $p\bar{p} \rightarrow \Lambda \Sigma^0 \bar{\Lambda} \bar{\Sigma}^0$. Another possibility would be to assume a mixture of $K$ and $K^*$ exchange as $K$ exchange gives a flatter angular distribution. Such an assumption could be tested easily from the asymptotic behaviour of the differential cross-sections with increasing energy. The amplitudes originating from the exchange of the $K$ meson decreases as $\frac{1}{s}$ relative to the amplitudes from $K^*$ exchange 12) and one should expect the weak peripheral events to vanish with increasing energy. This is not the case in the energy interval we study. As $K$ exchange also would introduce more parameters in our calculations we decided to disregard the weak peripheral events and assume that $K^*$ exchange was responsible for the marked peak in the forward direction and normalize the theoretical curves to the events in the interval $|t|_{\text{min}}$ to $|t| \sim 0.5 \text{ (GeV/c)}^2$.

The results of the calculation are shown in Fig. 5 and the agreement between theory and experiment is fairly good at low momentum transfer.

The energy dependence of the effective coupling constants is much weaker than for the $\Lambda\bar{\Lambda}$ production as can be seen from Fig. 7.

*) They also cannot be compared to the events in the reaction $p\bar{p} \rightarrow \Sigma^- \bar{\Sigma}^-$, which show a peak in the forward and in the backward direction.
We examine next the process $p\bar{p} \rightarrow \Lambda \overline{\Sigma}^0 + \Sigma^0 \overline{\Lambda}$, and determine the product of the effective coupling constants

$$G_{\Lambda \Sigma} = \frac{2}{\pi} \left( \frac{k_{N \Lambda}}{4 \pi} \right)^2 \left( \frac{k_{N \Sigma}}{4 \pi} \right)^2$$

by normalizing, as previously, the calculated curves to the experimental curves. The results are shown in Fig. 6. The angular distribution is reasonable fitted but we are nevertheless far from the excellent agreement of the $\Lambda\overline{\Lambda}$ case. The calculated curves are normalized to the total cross-section although there are some weak peripheral events.

Having determined already $\varepsilon_{K^*\Lambda N}^2$ and $\varepsilon_{K^*\Sigma^0 N}^2$ from the two first processes, this last process should not involve any new parameter if the elastic interaction in the final state is the same. We are thus in a position to check the one particle exchange model by comparing, for example, the value $\varepsilon_{K^*\Sigma^0 N}^2/4\pi$ obtained from experiment and its value obtained from $\varepsilon_{K^*\Lambda N}^2$ and $G_{\Lambda \Sigma}$ through (3.1); the following relation should hold

$$\frac{\varepsilon_{K^*\Sigma^0 N}^2}{4\pi} = \frac{\varepsilon_{K^*\Lambda N}^2}{4\pi} = \frac{G_{\Lambda \Sigma}}{4\pi}$$

The comparison is made in Table 1. We note that, for $P = 3$ GeV/c and equal elastic scattering parameters, we get incompatibility between the two values. However, the agreement is better at higher energies, thanks to the great uncertainty in the experimental results and to the fact that the three different cross-sections become less different. It is however easy in the framework of the model to make the experimental results compatible with the $K^*$ exchange, by assuming different elastic interactions in the final states than in the final states. As a specific example, we give in the last line of Table 1 the results of calculations at 3 GeV, assuming

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\[ 3 A(\Lambda \Lambda) = 3 A(\Lambda \Sigma) = A(p \bar{p}) = \frac{4}{3} A(\Sigma^+ \Sigma^+) \]

This affects very little the angular dependence compared to the first calculations. An assumption of this kind can also give the same energy dependence for the effective coupling constants \( \frac{g_{K^*N\Sigma}^2}{4\pi} \) and \( \frac{g_{K^*N\Lambda}^2}{4\pi} \). To try, however, to explain along these lines the whole energy variation of the coupling constants would most probably be meaningless. First, we would need to assume very great differences between the \( p \bar{p} \) and \( Y\bar{Y} \) elastic scattering. Also, as the cross-section of \( Y\bar{Y} \) production must fall at high energy, we would be forced to assume a strong shrinking of the elastic diffraction peak with increasing energy.

<table>
<thead>
<tr>
<th>( P_{\text{lab. in GeV/c}} )</th>
<th>( \frac{g_{K^*\Lambda N}^2}{4\pi} )</th>
<th>( g_{\Lambda \Sigma} )</th>
<th>( \frac{g_{\Lambda \Sigma}}{\frac{g_{K^*\Lambda N}^2}{4\pi}} )</th>
<th>( \frac{g_{K^*N\Sigma^+}^2}{4\pi} )</th>
</tr>
</thead>
<tbody>
<tr>
<td>3 (CERN)</td>
<td>3.7 ± 0.3</td>
<td>12.9 ± 2.1</td>
<td>3.5 ± 1</td>
<td>1.6 ± 0.15</td>
</tr>
<tr>
<td>3.6 (do)</td>
<td>2.8 ± 0.4</td>
<td>6.8 ± 1.9</td>
<td>2.5 ± 0.9</td>
<td>1.5 ± 0.25</td>
</tr>
<tr>
<td>5.7 (do)</td>
<td>1.4 ± 0.2</td>
<td>2.0 ± 0.4</td>
<td>1.5 ± 0.5</td>
<td>1.2 ± 0.15</td>
</tr>
<tr>
<td>3.25 (Yale)</td>
<td>3.1 ± 0.25</td>
<td>6.3 ± 1.2</td>
<td>2.1 ± 0.6</td>
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</tr>
<tr>
<td>3.69 (do)</td>
<td>3.0 ± 0.25</td>
<td>7.7 ± 1.4</td>
<td>2.6 ± 0.6</td>
<td></td>
</tr>
<tr>
<td>3 (CERN)</td>
<td>2.4 ± 0.25</td>
<td>5.6 ± 1.0</td>
<td>2.2 ± 0.5</td>
<td></td>
</tr>
</tbody>
</table>

Table 1: values of the effective coupling constants and the ratio \( g_{\Lambda \Sigma} / (\frac{g_{K^*\Lambda N}^2}{4\pi}) \)
The values of $\varepsilon^2_{K^{*}N}/4\pi$, $G_{\Lambda\Sigma}$ and $\varepsilon^2_{K^{*}\Sigma^{+}}/4\pi$ are determined at each energy by normalizing the calculated results to the experimental data. In the first five rows we have assumed equal shape of the different elastic diffraction peaks: $A(\Lambda^{\Lambda}) = A(\Lambda^{0}) = A(p\bar{p}) = A(\Sigma^{+}\Sigma^{+})$, in the last row we have used $3A(\Lambda^{\Lambda}) = 3A(\Lambda^{0}) = A(p\bar{p}) = \frac{1}{3}A(\Sigma^{+}\Sigma^{+})$. The error limits indicated are based on uncertainties in the experimental cross-sections only.

CONCLUSION

These calculations reflect both the satisfactory and unsatisfactory aspects of the absorption model: on one side a fair description of the angular dependence of the produced particles (except the weak peripheral events in the case of $\Sigma^{+}\Sigma^{+}$ production); on the other side, a wrong energy dependence of the cross-sections caused by the exchange of a particle with spin.

In the framework of the model, this last result can be understood as a consequence of the assumption that the effects of all the computing reaction channels can be represented by elastic scattering parameters only.

Regge pole theory would in principle give an easy solution to this difficulty, but in practice its predictive power is still very small. Thus, the absorption model and Regge pole model appear to have their strength on different aspects; Regge poles make it easy to get the energy dependence in the forward direction but are completely helpless in describing the variation with $t$, as the residue function of the poles are, until now, beyond anybody's computational capacity. The absorption model, on the other hand, gives an intuitive understanding of the angular dependence but suffers from the well-known difficulties associated with the exchange of elementary particles with spin.
ACKNOWLEDGEMENTS

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APPENDIX

THE HELICITY BORN AMPLITUDES

We show in Table 2 how the sixteen amplitudes are expressed from six of them:

<table>
<thead>
<tr>
<th>$\lambda_a$ $\lambda_c$</th>
<th>$++$</th>
<th>$+-$</th>
<th>$-+$</th>
<th>$--$</th>
</tr>
</thead>
<tbody>
<tr>
<td>$\lambda_b$ $\lambda_d$</td>
<td>$M_1$</td>
<td>$M_4$</td>
<td>$-M_3$</td>
<td>$M_2$</td>
</tr>
<tr>
<td>$++$</td>
<td>$M_4$</td>
<td>$-M_6$</td>
<td>$M_5$</td>
<td>$-M_3$</td>
</tr>
<tr>
<td>$+-$</td>
<td>$M_3$</td>
<td>$M_5$</td>
<td>$-M_6$</td>
<td>$M_4$</td>
</tr>
<tr>
<td>$-+$</td>
<td>$M_2$</td>
<td>$M_3$</td>
<td>$M_4$</td>
<td>$M_1$</td>
</tr>
</tbody>
</table>

Table 2: relations between the helicity amplitudes.

In order to write their complete expression it is useful to introduce the kinematical quantities

$$ f_{\pm} = \sqrt{(E + m_N)(E + m_Y)} \left( \frac{q}{m_N + E} \pm \frac{q'}{m_Y + E} \right) $$

(A.1)

$$ g_{\pm} = \sqrt{(E + m_N)(E + m_Y)} \left( 1 \pm \frac{qq'}{(m_N + E)(m_Y + E)} \right) $$

(A.2)

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The six independent amplitudes are now expressed in terms of $\Theta_\pm$, $\eta_{\pm}$, the production angle $\Theta$, and the quantities $G_{VV}$, $G_{4S}$ and $G_{SS}$, defined in (2.5). Note that $G_{SS}$ depends on $\Theta$.

\[(m^2-t)M_1 = G_{VV} \left( \Theta_+^2 + \sin^2 \frac{\Theta}{2} \Theta_+^2 \cos \frac{\Theta}{2} \eta_+^2 \right) + G_{4S} \left( 2 \cos \frac{\Theta}{2} \eta_+^2 + G_{SS} \cos \frac{\Theta}{2} \eta_+^2 \right) - (m^2-t)M_2 = G_{VV} \cos \frac{\Theta}{2} \left( \Theta_+^2 + \eta_+^2 \right) + 2G_{4S} \cos \frac{\Theta}{2} \eta_+^2 + G_{SS} \cos \frac{\Theta}{2} \eta_+^2 \right) - (m^2-t)M_3 = \cos \frac{\Theta}{2} \left( \Theta_+^2 + \eta_+^2 \right) + 2G_{4S} \cos \frac{\Theta}{2} \eta_+^2 + G_{SS} \cos \frac{\Theta}{2} \eta_+^2 \right) \]

\[(m^2-t)M_4 = G_{VV} \sin \frac{\Theta}{2} \cos \frac{\Theta}{2} \left( \Theta_+^2 + \eta_+^2 \right) - G_{4S} \sin \frac{\Theta}{2} \left( \eta_+^2 + \eta_-^2 \right) - G_{SS} \sin \frac{\Theta}{2} \cos \frac{\Theta}{2} \eta_+^2 \right) - (m^2-t)M_5 = \cos \frac{\Theta}{2} \left( \Theta_+^2 + \eta_+^2 \right) + 2G_{4S} \sin \frac{\Theta}{2} \eta_+^2 + G_{SS} \sin \frac{\Theta}{2} \eta_+^2 \right) \]

\[(m^2-t)M_6 = G_{VV} \left( \Theta_+^2 + \cos \frac{\Theta}{2} \Theta_+^2 \cos \frac{\Theta}{2} \eta_+^2 \right) + 2G_{4S} \sin \frac{\Theta}{2} \eta_+^2 + G_{SS} \sin \frac{\Theta}{2} \eta_+^2 \right) \]

For the partial wave expansion we split $G_{SS}$ (with $g_S = 0$) in two terms

\[G_{SS} = G_{SS}^0 - 4G_{SS}^1 \sin^2 \Theta \]

where

\[G_{SS}^0 = -2q_T(q_T + q_V) - \frac{q_T^2}{(m_T + m_Y)^2} (4E^2 + (q_T + q_V)^2) - \frac{(m_T + m_Y)^2}{m_T^2} q_V \]

\[G_{SS}^1 = \frac{q_T^2}{(m_T + m_Y)^2} q_V \]

(A.4)
and the expansion is easily done in terms of the $c$ functions defined in the text (2.7)

\[
M_1 = B_1 \sum_{j=0}^{\infty} c^{(2)}_{j=0} k(j) (2j+1) \\
M_2 = B_2 \sum_{j=0}^{\infty} c^{(2)}_{j=0} k(j) \\
M_3 = B_3 \sum_{j=0}^{\infty} c^{(2)}_{j=0} k(j) \\
M_4 = B_4 \sum_{j=0}^{\infty} c^{(2)}_{j=0} k(j) \\
M_5 = B_5 \sum_{j=0}^{\infty} c^{(2)}_{j=0} k(j) \\
M_6 = B_6 \sum_{j=0}^{\infty} c^{(2)}_{j=0} k(j)
\]

In the Born approximation all the $k(j)$'s are identical equal to 1, in the absorption model they are given by (2.9) in the text.

The expressions for the $A$'s and the $B$'s are given by $G_{VV}$, $G_{4S}$, $G_{SS}^{0}$ and $G_{SS}^{1}$ as follows

\[
B_1 = \frac{4}{4q^2} G_{VV} (\gamma_{+}^{2} - \gamma_{-}^{2}) + 2 G_{4S} (\gamma_{+}^{2} - \gamma_{-}^{2}) - 2 G_{SS}^{0} \gamma_{-}^{2} + G_{SS}^{1} \gamma_{-}^{2} \\
B_2 = \frac{1}{2} B_2 = \frac{G_{SS}^{0}}{2q^2} \gamma_{-}^{2}
\]
\[ B_3 = B_4 = \frac{12}{2qq'} G_{ss}^4 \eta \gamma \eta \]

\[ B_5 = -2 B_6 = \frac{G_{ss}^4}{q q'} \gamma^2 \]

\[ B_8 = \frac{1}{4qq'} \left\{ G_{uv} (\eta^2 - \xi^2) + 2 G_{qs} \gamma \eta \gamma + G_{ss}^6 \gamma^2 + 2 G_{ss}^4 \gamma^2 \xi^2 \right\} \]

\[ A_1 = \frac{1}{2qq} \left\{ \frac{1+2}{z} (G_{uv} (\eta^2 - \xi^2) + 2 G_{qs} \gamma \eta \gamma + G_{ss}^6 \gamma^2) + G_{ss}^4 \gamma^2 (1-z)^2 + 2 G_{uv} \xi^2 \right\} \]

\[ A_2 = \frac{1}{2qq} \left\{ G_{uv} (\gamma^2 + \eta^2) + 2 G_{qs} \gamma \eta - G_{ss}^6 \gamma^2 + 2 G_{ss}^4 \gamma^2 (1-z) \right\} \]

\[ A_3 = \frac{1}{2qq} \left\{ G_{uv} (\gamma^2 + \eta^2) + 2 G_{qs} (\gamma^2 + \eta^2) + G_{ss}^6 \gamma^2 + 2 G_{ss}^4 \gamma^2 (1-z) \right\} \]

\[ A_4 = \frac{1}{2qq} \left\{ G_{uv} (\xi^2 + \eta^2) + 2 G_{qs} \gamma \eta - G_{ss}^6 \gamma^2 + 2 G_{ss}^4 \gamma^2 (1-z) \right\} \]

\[ A_5 = \frac{1}{2qq} \left\{ G_{uv} (\xi^2 + \eta^2) + 2 G_{qs} \gamma \eta - G_{ss}^6 \gamma^2 + 2 G_{ss}^4 \gamma^2 (1-z) \right\} \]

\[ A_6 = \frac{1}{2qq} \left\{ \frac{1-z}{2} \left[ G_{uv} (\eta^2 - \xi^2) + 2 G_{qs} \gamma \eta - G_{ss}^6 \gamma^2 \right] + G_{ss}^4 \gamma^2 (1-z)^2 + 2 G_{uv} \xi^2 \right\} \]
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FIGURE CAPTIONS

Figure 1  The one-particle exchange diagram.

Figure 2  Experimental data for the total cross-sections, CERN 1) data are marked with ○, YALE 2) data with □.

Figure 3  The differential cross-section for $\Lambda\bar{\Lambda}$ production, at $P = 3 \text{ GeV/c}$ and $3.6 \text{ GeV/c}$ (CERN data).

Figure 4  The differential cross-section for $\Lambda'\bar{\Lambda}'$ production, at $P = 5.7 \text{ GeV/c}$ and $7 \text{ GeV/c}$ (CERN data at $5.7 \text{ GeV/c}$, unpublished), (YALE data at $7 \text{ GeV/c}$, Ref. 2)).

Figure 5  The differential cross-section for $\Sigma^+\bar{\Sigma}^+$ production (CERN data).

Figure 6  The differential cross-section for $(\Lambda\Sigma^0 + \Sigma^0\Lambda)$ production (CERN data).

Figure 7  The effective coupling constants as a function of energy plotted in logarithmic scale:

- ○ denotes $\varepsilon_{K^*N\Lambda}^2/4\pi$ with $A(\Lambda\bar{\Lambda}) = A(p\bar{p})$
- ○ denotes $\varepsilon_{K^*N\Lambda}^2/4\pi$ with $3A(\Lambda\bar{\Lambda}) = A(p\bar{p})$
- □ denotes $\varepsilon_{K^*N\Sigma^0}^2/4\pi$ with $A(\Sigma^+\Sigma^+\bar{\Lambda}) = A(p\bar{p})$
- □ denotes $\varepsilon_{K^*N\Sigma^0}^2/4\pi$ with $A(\Sigma^+\Sigma^+\bar{\Lambda}) = 3A(p\bar{p})$
FIG. 1
FIG. 3

$d\sigma/d(\cos \theta)$ in $\mu$b/$\cos \theta$

$P = 3$ GeV/c

$P = 3.6$ GeV/c