# K* (890) PHOTOPRODUCTION IN THE REGGE POLE MODEL* 

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#### Abstract

In this paper the Regge pole model is applied to $K^{*}(890)$ photoproduction. Conspiracy relations, kinematic constraints, factorization, and the question of when to set $m_{\gamma}=0$ are discussed. Predictions are obtained, in the limit of large $s$ and small $t$, for $\frac{\mathrm{d} \sigma}{\mathrm{d} \Omega}$, the density matrix, and the $\mathrm{K}^{*}$ decay angular distribution.


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[^0]
## I. INTRODUCTION

Recently much interest has been shown in the application of Regge poles to high energy processes. ${ }^{1} \mathrm{~K}^{*}(890)$ photoproduction is an interesting reaction to consider here; the particles all have spin, and the masses are all unequal.

In this paper we apply the Regge pole model to the process

$$
\begin{equation*}
\gamma+\mathrm{p} \longrightarrow \mathrm{~K}^{*}+\mathrm{Y} \tag{1}
\end{equation*}
$$

where $\mathrm{K}^{*}$ is the $1^{-}$meson at 890 MeV , and Y is either a $\Lambda$ or a $\sum$ hyperon. Kinematic singularities are separated out of the helicity amplitudes, and the questions of conspiracy relations, kinematic constraints, factorization, and when to set $\mathrm{m}_{\gamma}=0$ are investigated. We investigate the values of t where constraints may arise and discover that the leading amplitudes for the process (1) are not involved in any constraint relations at $t=0$ or at $t=\left(m_{p}-m_{Y}\right)^{2}$. We also discuss why the crossing matrix should yield the same conspiracy relations as those obtained from invariant amplitudes. Predictions for the large $s$ and small $t$ behavior of $\frac{d \sigma}{d \Omega}$, the density matrix $\rho$, and the $K^{*}$ decay angular distribution are obtained; no data is as yet available for comparison with these predictions.

The plan of this paper is as follows: In Section II we define our kinematic-singularity-free amplitudes and find the dominant contributions in the limit of large $s$ and small $t$. Section III discusses possible modifications due to conspiracy relations, kinematic constraints, factorization, and $\mathrm{m}_{\gamma} \longrightarrow 0$. In Section IV we obtain our predictions for $\frac{\mathrm{d} \sigma}{\mathrm{d} \Omega}$, the density matrix, and the $\mathrm{K}^{*}$ angular distribution.

## II. $t$ CHANNEL HELICITY AMPLITUDES

In this section we shall investigate the $t$ channel helicity amplitudes associated with $\mathrm{K}^{*}(890)$ photoproduction. We shall construct t channel helicity amplitudes that are free of kinematical singularities; these will then be Reggeized. In
separating out the kinematic singularities we shall follow the method of Wang ${ }^{2}$; possible modifications will be discussed in Section III.

Each $t$ channel helicity amplitude $f_{c d ; a b}^{t}$ contains factors of $\sin \left(\theta_{t} / 2\right)$ and $\cos \left(\theta_{\mathrm{t}} / 2\right)$ arising from a partial wave expansion in terms of $d$ functions; these are separated out according to

$$
\begin{equation*}
\mathrm{f}_{\mathrm{cd} ; \mathrm{ab}}^{\mathrm{t}}=\left(\sin \frac{\theta_{\mathrm{t}}}{2}\right)^{|\lambda-\mu|}\left(\cos \frac{\theta_{\mathrm{t}}}{2}\right)^{|\lambda+\mu|} \overline{\mathrm{f}}_{\mathrm{cd} ; \mathrm{ab}}^{\mathrm{t}} \tag{2}
\end{equation*}
$$

where $\lambda=\mathrm{a}-\mathrm{b}$ and $\mu=\mathrm{c}-\mathrm{d} . \quad \widetilde{\mathbf{f}}^{\mathrm{t} \pm}$ is then defined by

$$
\begin{equation*}
\overline{\mathrm{f}}_{\mathrm{cd} ; \mathrm{ab}}^{\mathrm{t}} \pm \overline{\mathrm{f}}_{-\mathrm{c}-\mathrm{d} ; \mathrm{ab}}^{\mathrm{t}}=\overline{\mathrm{K}}_{\mathrm{cd} ; \mathrm{ab}}^{\mathrm{t}}(\mathrm{t}){\underset{\mathrm{f}}{\mathrm{~cd} ; \mathrm{ab}}}_{\mathrm{t} \pm}^{(\mathrm{t}, \mathrm{~s})} \tag{3}
\end{equation*}
$$

where $\overline{\mathrm{K}}$ contains the kinematic singularities. $\widetilde{\mathrm{f}}^{\mathrm{t}}$ is the kinematic-singularityfree amplitude which is Reggized according to

$$
\begin{equation*}
\widetilde{\mathrm{f}}_{\mathrm{cd} ; a b}^{\mathrm{t}}(\mathrm{t}, \mathrm{~s}) \longrightarrow \gamma(\mathrm{t})\left(\frac{1 \pm \mathrm{e}^{-\mathrm{i} \pi \alpha(\mathrm{t})}}{\sin \pi \alpha(\mathrm{t})}\right)\left(\frac{\mathrm{s}}{\mathrm{~s}_{0}}\right)^{\alpha(\mathrm{t})-\mathrm{M}} \tag{4}
\end{equation*}
$$

where $M=\max \{|\lambda|,|\mu|\}$. The factor $M$ would absent in the spinless case; in the case with spin it arises because some of the powers of $s$ are absorbed by separating out the $\sin \left(\theta_{\mathfrak{t}} / 2\right)$ and $\cos \left(\theta_{\mathrm{t}} / 2\right)$ factors before Reggeization. The rest of the powers of $s$ are assumed to contribute full strength in our unequal mass case (i.e., we are assuming the action of daughter trajectories ${ }^{3}$ when writing $\left.\mathrm{s}^{\alpha(\mathrm{t})-\mathrm{M}}\right)$.

We note that for small $t$ and a given $\alpha(t)$, the highest power of $s$ occurs when $M=0$. (The $\sin \theta_{\mathrm{t}} / 2$ and $\cos \theta_{\mathrm{t}} / 2$ factors do not contribute any powers of s herc, since we arc dealing with an unequal mass case $-\mathrm{m}_{\gamma} \neq \mathrm{m}_{\mathrm{K}^{*}} \cdot$ ) Applying parity conservation ${ }^{4}$ to the $\gamma \mathrm{K}^{*}$-Reggion vertex and cxpanding in partial waves, we find that exchange of natural parity trajectorics $\left(P=(-1)^{\mathrm{J}}\right)$ can contribute to $\widetilde{\mathrm{f}}_{\lambda \neq \mu=0}^{\mathrm{t}+}$, but not to $\widetilde{\mathrm{F}}_{\lambda \neq \mu=0}^{\mathrm{t}-}$. Thus the amplitudes we shall need are those
$\widetilde{\mathbf{f}}^{\mathbf{t +}}$ amplitudes having $\lambda=\mu=0 .{ }^{5}$ Using parity ${ }^{4}$ to reduce the number of independent amplitudes, we are left with the following amplitudes at large $s$ and small t :

$$
\begin{equation*}
\mathrm{f}_{11 ; \frac{1}{2} \frac{1}{2}}^{\mathrm{t}}=\mathrm{f}_{-1-1 ;-\frac{1}{2}-\frac{1}{2}}^{\mathrm{t}} \cong \mathrm{f}_{-1-1 ; \frac{1}{2} \frac{1}{2}}^{\mathrm{t}}=\mathrm{f}_{11 ;-\frac{1}{2}-\frac{1}{2}}^{\mathrm{t}} \cong \frac{1}{2} \overline{\mathrm{~K}}_{11 ; \frac{1}{2} \frac{1}{2}}^{\mathrm{t}} \widetilde{\mathrm{f}}_{11 ; \frac{1}{2} \frac{1}{2}}^{\mathrm{t}+} \tag{5}
\end{equation*}
$$

where

$$
\begin{equation*}
\widetilde{\mathrm{f}}_{11 ; \frac{1}{2} \frac{1}{2}}^{\mathrm{t}+} \longrightarrow \gamma(\mathrm{t})\left(\frac{1 \pm \mathrm{e}^{-\mathrm{i} \pi \alpha(\mathrm{t})}}{\sin \pi \alpha(\mathrm{t})}\right)\left(\frac{\mathrm{s}}{\mathrm{~s}_{0}}\right)^{\alpha(\mathrm{t})} \tag{6}
\end{equation*}
$$

and $\overline{\mathrm{K}}_{11 ; 22}^{+}{ }^{(\mathrm{t})}$ is a kinematical factor. Using the prescription given by Wang for the unequal mass case (and setting $\mathrm{m}_{\gamma} \rightarrow 0$ in this result), we find that

$$
\begin{equation*}
\overline{\mathrm{K}}_{11 ; \frac{1}{2} \frac{1}{2}}^{+}(\mathrm{t})=\left[\mathrm{t}-\left(\mathrm{m}_{\mathrm{p}}+\mathrm{m}_{\mathrm{Y}}\right)^{2}\right]^{-1 / 2}\left[\mathrm{t}-\mathrm{m}_{\mathrm{K}^{*}}^{2}\right]^{-2} \tag{7}
\end{equation*}
$$

Actually, some caution is ncedcd with respect to this form of $\overline{\mathrm{K}}^{+}$, and before proceeding further, we shall discuss possible modifications of the factor $\overline{\mathrm{K}}^{+}$due to conspiracy rclations, kinematic constraints, factorization, and the question of when to set $m_{\gamma}=0$.
III. CONSPIRACY RELATIONS KINEMATIC CONSTRAINTS, FACTORIZATION, AND $\mathrm{M}_{\gamma} \rightarrow 0$.
There are several points to be discussed in connection with the factor $\overline{\mathrm{K}}_{11 ; \frac{1}{2} \frac{1}{2}}^{+}(\mathrm{t})$. The general form of $\overline{\mathrm{K}}$ has been derived by Wang, ${ }^{2}$ who examined the singularities in the crossing matrix relating $s$ and $t$ channel helicity amplitudes. In addition, however, the questions of conspiracy relations, kinematic constraints, factorization, and how to treat $\mathrm{m}_{\gamma} \rightarrow 0$ may arise. Consideration of these points could lead to a modified $\overline{\mathrm{K}}$.

One can ask at which points to expect special conditions on the $t$ channel helicity amplitudes. In $\overline{\mathrm{N}} \mathrm{N}$ scattering angular momentum conservation applied at $\cos \theta_{S}= \pm 1$ leads to conspiracy relations, ${ }^{6}$ but in the unequal mass case (i.e. , $\mathrm{m}_{\gamma} \nLeftarrow \mathrm{m}_{\mathrm{K}^{*}}$ ) there are no such relations at $\cos \theta_{\mathrm{s}}= \pm 1$. (This has been shown by Högaasen and Salin; ${ }^{7}$ the proof depends on the fact that for unequal masses, $\cos \theta_{\mathrm{S}}= \pm 1$ implies $\cos \theta_{\mathrm{t}}= \pm 1$. Angular momentum conservation then says that each $f_{\lambda \neq \mu}^{\mathrm{t}}$ must vanish at $\cos \theta_{\mathrm{S}}= \pm 1$, and no relations between $\mathrm{f}^{\mathrm{t}} \mathrm{s}$ arise.)

Conspiracy relations can still arise at $t=0$ when all the masses are unequal, but no such relations occur for $\lambda=0$ or $\mu=0,{ }^{8}$ which is the case of interest here.

In general we might also expect special relations between helicity amplitudes to arise at those points where the helicity becomes undefined. ${ }^{9}$ One can define the helicity four-vector $n_{3}\left(p_{i}\right)$ for a two-particle state by the conditions $n_{3} \cdot n_{3}=-1$ and (in the c.m. system) $\vec{n}_{3} \cdot \vec{p}_{i}>00^{9}$

$$
n_{3}\left(p_{i}\right)=-\left\{\frac{m_{i}^{2} P-\left(p_{i} \cdot P\right) p_{i}}{\frac{m_{i}}{2}\left[t-\left(m_{1}+m_{2}\right)^{2}\right]^{1 / 2}\left[t-\left(m_{1}-m_{2}\right)^{2}\right]^{1 / 2}}\right\}
$$

where $P=p_{1}+p_{2}$. It is evident that troubles arise when. $t$ is at a threshold or pseudothreshold $t=\left(m_{1} \pm m_{2}\right)^{2}$, and constraint conditions can occur between $t$ channel amplitudes at precisely these points. Another way of seeing that constraint conditions arise at $t=\left(m_{i} \pm m_{j}\right)^{2}$ has been discussed by Jackson and Hite, ${ }^{10}$ who note that in a special basis system certain amplitudes vanish atprecisely these points.

Having thus discussed where one can expect constraint conditions, we next turn to the question of constructing them. One way of deriving constraint conditions is to express the invariant scalar amplitudes for the process in terms of linear combinations of $t$ channel helicity amplitudes. Since the scalar amplitudes have no poles in $t$, one then derives certain conditions on linear combinations of t channel helicity amplitudes. These are the conspiracy relations. The point to be noted is that the only property of the scalar amplitudes that is used is that
they have no poles in $t$. Thus one could have started with any other complete set of $s$ channel amplitudes having no poles in $t$, and the results would have been exactly the same. Hence it is equally valid to start with $f^{s}$ amplitudes. Since these are related to the $f^{t^{t}} \mathrm{~s}$ by crossing, one can thus obtain the conspiracy relations by examining the crossing matrix at the values of $t$ in question. This approach has been investigated in detail by Cohen-Tannoudji et al. ${ }^{9}$

We now turn to explicit construction of the desired constraint relations for $\mathrm{K}^{*}$ photoproduction. We need those relations which involve $\lambda=\mu=0$ helicity amplitudes (these are the relevant amplitudes for large s, as noted in Section II); the point of interest is $t=t_{0}=\left(m_{p}-m_{Y}\right)^{2}$. The other three threshold or psuedothreshold points for this reaction involve much larger values of $t$ and are thus not needed for a study of the behavior at small t. Following the method involving the crossing matrix ${ }^{9}$ as illustrated by Högaasen and Salin, ${ }^{6}$ we discover that there are indeed constraint relations between several $t$ channel helicity amplitudes at $t=t_{0}$, but none of these relations involves $\overline{\mathrm{f}}_{\mathrm{cc} ; \mathrm{aa}}^{\mathrm{t}}+\overline{\mathrm{f}}_{-\mathrm{c}-\mathrm{c} ; \mathrm{aa}}^{\mathrm{t}} \quad$ amplitudes. I. e., those amplitudes which contain the leading $s$ behavior (as discussed above) are not involved in any conspiracy or constraint relations.

Thus result can be partially understood in the following way: The $\bar{f}_{c c ; a a}^{t}+\bar{f}_{-c-c ; a a}^{t} s$ cannot be linearly related to other $\overline{\mathrm{f}}^{\mathrm{t}_{\mathrm{s}}} \mathrm{s}$ whose leading s behavior is also governed by natural parity $\left(P=(-1)^{J}\right)$ exchanges at $t=t_{0}$, since the leading powers of $s$ would be different. On the other hand, $\overline{\mathrm{f}}_{\mathrm{cc} ; \mathrm{aa}}^{\mathrm{t}}+\overline{\mathrm{f}}_{-\mathrm{c}-\mathrm{c} ; \mathrm{aa}}^{\mathrm{t}}$ can not be related to a polynomial in $s$ times otheri $\bar{f}^{t}$, since none of the $\overline{\mathrm{f}}^{\mathrm{t}} \mathrm{s}$ has any kinematic singularities in $s$.

Thus in finding the kinematical factor $\overline{\mathrm{K}}_{\mathrm{cc}}^{+}$;aa we can completely avoid the question of possible extra factors due to conspiracy or constraint relations. We next note that factorization of residue functions will also not yield any new
information. The residue factors $\overline{\mathrm{K}}$ factor automatically at thresholds and pseudothresholds, and for our unequal mass $\lambda=\mu=0$ helicity amplitudes there are no factors of $t$ in $\overline{\mathrm{K}}$ for the relevant processes, so factorization with respect to these pieces is automatically satisfied.

To find the kinematical factors, we thus gain no new information from constraint conditions or from factorization. The only remaining uncertainty in finding $\overline{\mathrm{K}}_{\mathrm{cc} ; \mathrm{aa}}^{+}$arises from the question when to put $\mathrm{m}_{\gamma}=0$. One could put $\mathrm{m}_{\gamma}=0$ in Wang's general prescription ${ }^{2}$ for $\bar{K}_{c c ; a a}^{+}$, or one could put $m_{\gamma}=0$ in the crossing matrix and derive a modified prescription for $\overline{\mathrm{K}}_{\mathrm{cc} ; \mathrm{aa}}^{+} \cdot{ }^{11}$ The end result for the two cases can in general differ in the net power to which ( $t-m_{K^{*}}^{2}$ ) should be raised. Since for small $t$ this factor is smooth, the exact power need not concern us, and we will simply use Wang's prescription and set $\mathrm{m}_{\gamma} \longrightarrow 0$ at the end. We thus obtain the result

$$
\begin{equation*}
\overline{\mathrm{K}}_{11 ; \frac{1}{2} \frac{1}{2}}^{+}(\mathrm{t})=\left[\mathrm{t}-\left(\mathrm{m}_{\mathrm{p}}+\mathrm{M}_{\mathrm{Y}^{\prime}}\right)^{2}\right]^{-1 / 2}\left[\mathrm{t}-\mathrm{m}_{\mathrm{K}^{*}}^{2}\right]^{-2} \tag{7}
\end{equation*}
$$

This result for the kinematic factor $\overline{\mathrm{K}}$ can be checked by means of simple angular momentum and parity arguments.$^{12}$ As an example of the method, consider the point $t=\left(m_{p}+M_{Y}\right)^{2}$. Setting $\vec{l}$ (orbital angular momentum) for the $N \bar{Y}$ system equal to zero, the possible $N \bar{Y}$ states have $J^{P}=0^{-}$or $1^{-}$. Since the $N \bar{Y}$ system is coupled to a $P=(-1)^{J}$ Reggion in our model, $J P$ is restricted to $1^{-}$. Now expand the helicity amplitudes in terms of partial waves:

$$
\overline{\mathrm{f}}_{11 ; \frac{1}{2} \frac{1}{\mathrm{t}}}^{\mathrm{t}}+\overline{\mathrm{f}}_{-1-1 ; \frac{1}{2} \frac{1}{2}}^{\mathrm{t}}=\sum_{\mathrm{J}} \mathrm{~F}_{11 ; \frac{1}{2} \frac{1}{\mathrm{~J}}}^{\mathrm{J}} \mathrm{~d}_{00}^{\mathrm{J}}\left(\cos \theta_{\mathrm{t}}\right)
$$

The $F^{J_{t}}$ have the threshold behavior $q_{N \bar{Y}}^{\ell}$, and $\cos \theta_{t}$ is proportional to $1 / q_{N} \bar{Y}$.

Hence we deduce the behavior:

$$
F^{J} d_{00}^{J}\left(\cos \theta_{t}\right) \propto q_{N \bar{Y}}^{\ell}\left(1 / q_{N \bar{Y}}\right)^{J}=\left(q_{N} \bar{Y}\right)^{-1},
$$

i.e., near $t=\left(m_{p}+m_{Y}\right)^{2}$ the kinematic factor $\bar{K}^{+}(t)$ goes as $q_{N Y}{ }^{-1} \propto\left[t-\left(m_{p}+m_{Y}\right)^{2}\right]^{-1 / 2}$, in agreement with (7). For $t$ near the other threshold or pseudothreshold points, analogous arguments go through (the $\bar{Y}$ is treated as having posilive parity at $\left.t=\left(m_{p}-m_{Y}\right)^{2}\right)$, and one obtains exactly the Wang kinematic factor (7).
IV. CROSS SECTION, DENSITY MATRIX, AND $K^{*}$ DECAY ANGULAR DISTRIBUTION

The differential cross section in the $c, m$. frame can be written

$$
\frac{\mathrm{d} \sigma}{\mathrm{~d} \Omega}=\frac{\mathrm{p}_{\mathrm{f}}}{4 \pi^{2} \mathrm{sp}_{\mathrm{i}}} \frac{1}{4} \sum\left|\mathrm{f}^{\mathrm{s}}\right|^{2}
$$

where the sum goes over the s-channel helicity amplitudes. Orthogonality of the crossing matrix ${ }^{13}$ then gives

$$
\begin{equation*}
\frac{\mathrm{d} \sigma}{\mathrm{~d} \Omega}=\frac{\mathrm{p}_{\mathrm{f}}}{4 \pi^{2} \mathrm{~s} \mathrm{p}_{\mathrm{i}}} \frac{1}{4} \sum\left|\mathrm{f}^{\mathrm{t}}\right|^{2} \tag{8}
\end{equation*}
$$

where the sum is now over all $t$-channel helicity amplitudes. We take the limit of large $s$ and use the results (5), (6), and (7) of the preceding sections. Thus

$$
\begin{align*}
&\left.\frac{d \sigma}{d \Omega}\right|_{c . m} . \begin{array}{l}
\text { small } \mathrm{t}
\end{array} \frac{\mathrm{p}_{\mathrm{f}}}{16 \pi^{2} \mathrm{sp}_{\mathrm{i}}}\left|\mathrm{t}-\left(\mathrm{m}_{\mathrm{p}}+\mathrm{m}_{\mathrm{Y}}\right)^{2}\right|^{-1}\left[\mathrm{t}-\mathrm{m}_{\mathrm{K}^{*}}^{2}\right]^{-4} \\
& \times\left|\sum_{\mathrm{i}} \gamma_{\mathrm{i}}(\mathrm{t}) \frac{1 \pm \mathrm{e}^{-\mathrm{i} \pi \alpha_{i}(\mathrm{t})}}{\sin \pi \alpha_{i}(\mathrm{t})}\left(\frac{\mathrm{s}}{\mathrm{~s}_{0}}\right)^{\alpha_{i}(\mathrm{t})}\right|^{2} \tag{9}
\end{align*}
$$

The sum is taken over the $\mathrm{K}^{*}\left(1^{-} ; 890\right)$, and $\mathrm{K}_{\mathrm{V}}\left(2^{+} ; 1420\right)$ trajectories. ${ }^{14}$
We make the following choices for the residue functions and Regge trajectories:
$\gamma_{1-}(t)$ is assumed roughly constant, while $\gamma_{2^{+}}(t)$ is put proportional to $\alpha_{2^{+}}(t)$ in order to cancel the pole in

$$
\frac{1+e^{-i \pi \alpha} 2^{+(t)}}{\sin \pi \alpha 2^{+}}(t) \quad
$$

at $\alpha_{2^{+}}{ }^{(t)}=0$ (Chew ghost-killing mechanism ${ }^{15}$ ). Thus we set

$$
\begin{align*}
& \gamma_{1^{-}}(t)=\widetilde{\gamma}_{1^{-}}(t)  \tag{10}\\
& \gamma_{2^{+}}(t)=\widetilde{\gamma}_{2^{+}}(t) \alpha_{2^{+}}(t)
\end{align*}
$$

where the $\tilde{\gamma}^{\prime} \mathrm{s}$ are assumed slowly varying in t . The $\mathrm{K}^{*}\left(1^{-}\right)$and $\mathrm{K}^{*}\left(2^{+}\right)$ trajectory functions are taken parallel to those of the $\rho$ and A2. We take ${ }^{16}$ $\alpha_{\rho}(\mathrm{t}) \cong \mathrm{t}+.57$ and $\alpha_{A 2}{ }^{(\mathrm{t})} \cong \mathrm{t}+.35$; thus

$$
\begin{align*}
& \alpha_{1^{-}}(t) \cong t+.37  \tag{11}\\
& \alpha_{2^{+}}(t) \cong t+.02
\end{align*}
$$

To determine the $\tilde{\gamma}_{i}^{\prime}$ s appearing in the residue functions $\gamma_{i}$, we first examine A2 and $\rho$ exchange in $\rho^{\circ}$ photoproduction; then we obtain the corresponding residues in $K^{*}$ photoproduction.

Since a photon does not couple to two $\rho^{\circ}$ 's by $C$ invariance, only the A2 contribution need be studied. Maheshwari ${ }^{17}$ has used universality and vector dominance to estimate the A 2 contribution to $\rho^{\circ}$ photoproduction. Evaluating his results at $t \cong 0$, we find that

$$
\tilde{\gamma}_{\mathrm{A} 2} \cong-\frac{\sqrt{\alpha}}{2} \quad \quad \gamma \mathrm{p} \rightarrow \rho^{\mathrm{o}} \mathrm{p}
$$

Thus for $\rho^{\circ}$ photoproduction $\tilde{\gamma}_{\rho}=0$ and $\tilde{\gamma}_{\mathrm{A} 2} \cong-\frac{\sqrt{\alpha}}{2}$. Using universality ${ }^{18}$ to rclate $\tilde{\gamma}_{\rho}$ and $\tilde{\gamma}_{\mathrm{A} 2}$ to $\gamma_{1^{-}}$and $\gamma_{2^{+}}$for $\mathrm{K}^{*}$ photoproduction, we obtain

$$
\begin{array}{ll}
\tilde{\gamma}_{1-}=0 & \\
\tilde{\gamma}_{2^{+}}=\frac{-\sqrt{2 \alpha}}{4} & , \gamma \mathrm{p} \rightarrow \mathrm{~K}^{*} \circ \sum^{+}  \tag{12}\\
\tilde{\gamma}_{2^{+}}=\frac{\sqrt{\alpha}}{4} & \gamma \mathrm{p} \rightarrow \mathrm{~K}^{*^{+}} \sum^{0} \\
\widetilde{\gamma}_{2^{+}}=\frac{-\sqrt{\alpha}}{4 \sqrt{3}} & \gamma \mathrm{p} \rightarrow \mathrm{~K}^{*^{+}} \Lambda^{\circ} . \\
& -y-
\end{array}
$$

We note that $\mathrm{g}_{\mathrm{K}_{1^{-}}^{*} \Sigma \mathrm{~N}}$ and $\mathrm{g}_{\mathrm{K}_{1}^{*}-\Lambda \mathrm{N}}$ are small ${ }^{19}$, so that $\tilde{\gamma}_{1^{-}}$is probably small even though $\mathrm{K}^{*}$ exchange is not prohibited by C invariance.

Thus we obtain the prediction (using (9), (10), (11), and (12) and setting $\left.\mathrm{s}_{\mathrm{o}} \cong 1(\mathrm{BeV})^{2}\right)$ :
$\left.\frac{\mathrm{d} \sigma}{\mathrm{d} \Omega}\right|_{\text {c. m. small }} \xrightarrow{\text { large } \mathrm{s}} \frac{\mathrm{p}_{\mathrm{f}}|\tilde{\gamma}|^{2} \alpha^{2}(\mathrm{t})\left|1+\mathrm{e}^{-\mathrm{i} \pi \alpha(\mathrm{t})}\right|^{2} \mathrm{~s}^{2 \alpha(\mathrm{t})}}{16 \pi^{2} \mathrm{~s} \mathrm{p}_{\mathrm{i}}\left|\mathrm{t}-\left(\mathrm{m}_{\mathrm{p}}+\mathrm{m}_{\mathrm{Y}}\right)^{2}\right|\left[\mathrm{t}-\mathrm{m}_{\mathrm{K}^{*}}^{2}\right]^{4} \sin ^{2} \pi \alpha(\mathrm{t})}$
where

$$
\alpha(\mathrm{t})=\mathrm{t}+.02
$$

and

$$
\begin{aligned}
\tilde{\gamma}(t) & =-\frac{\sqrt{2 \alpha}}{4} & & \gamma p \rightarrow \mathrm{~K}^{* o} \sum^{+} \\
& =\frac{\sqrt{\alpha}}{4} & & \gamma \mathrm{p} \rightarrow \mathrm{~K}^{*+} \sum^{o} \\
& =\frac{-\sqrt{\alpha}}{4 \sqrt{3}} & & \gamma \mathrm{p} \rightarrow \mathrm{~K}^{*+} \Lambda^{o} .
\end{aligned}
$$

The density matrix may also be found in our formalism. The density matrix can be expressed in terms of $t$ channel helicity amplitudes ${ }^{20}$ (we go to the $\mathrm{K}^{*}$ rest frame; the $z$ axis is taken parallel to the incident photon momentum as seen in this frame):

$$
\begin{equation*}
\rho_{m m^{\prime}}=\frac{\sum_{a, b, d} f_{m a ; d b}^{t^{*}} f_{m^{\prime} a ; d b}^{t}}{\sum_{a, b, c, d}\left|f_{c a ; d b}^{t}\right|^{2}} \tag{14}
\end{equation*}
$$

Inserting the result (5) into this expression, we obtain the predictions (for large $s$ and small t ):

$$
\begin{equation*}
\rho_{11}=\rho_{-1-1}=\frac{1}{2} ; \quad \rho_{00}=\rho_{\mathrm{m} \neq \mathrm{m}^{\prime}}=0 \tag{15}
\end{equation*}
$$

The $\mathrm{K}^{*}$ decay angular distribution has been written in terms of $\rho_{\mathrm{mm}}$, by Gottfried and Jackson. ${ }^{20,21}$

$$
\begin{align*}
\mathrm{W}(\theta, \phi)=\frac{3}{4 \pi}\left\{\rho_{00} \cos ^{2} \theta\right. & +\rho_{11} \sin ^{2} \theta-\rho_{1-1} \sin ^{2} \theta \cos 2 \phi \\
& \left.-\sqrt{2} \operatorname{Re} \rho_{10} \sin 2 \theta \cos \phi\right\} \tag{16}
\end{align*}
$$

Hence we directly obtain the prediction (for large $s$ and small $t$ ):

$$
\begin{equation*}
\mathrm{W}(\theta, \phi)=\frac{3}{8 \pi} \sin ^{2} \theta \tag{17}
\end{equation*}
$$

where $\theta$ is the angle made with the $z$ axis in the frame described above.
The only data on $\mathrm{K}^{*}$ photoproduction is an upper limit of 0.1 to $0.05 \mu \mathrm{~b}^{22}$ on the cross section for $\gamma+\mathrm{p} \longrightarrow \mathrm{K}^{\mathrm{O}^{*}}+\Sigma^{+}$, so no comparison with experiment can be made at the present time.

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