DIFFRACTION DISSOCIATION OF MESONS IN THE POMERON CURRENT MODEL

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Diffraction dissociation of pions and kaons is discussed in terms of the pomeron current model. Pomeron exchange is represented in this model by an effective current-current interaction, where the pomeron current V_μ has two components $V_\mu = V_{1\mu}+V_{2\mu}$, the first being central and conserved and the second peripheral and nonconserved. $V_{1\mu}$ and $V_{2\mu}$ correspond to interactions of sea and valence quarks respectively at the diffractive vertex. We construct explicit amplitudes for $N\pi \to N\pi^*$, $NK \to NK^*$ (π^*, K^* are resonances or partial waves of multiparticle states with any spin and parity). Two form factors (one form factor) determine the transition vertices $\pi \to \pi^*$, $K \to K^*$ to unnatural parity (natural parity) states π^* and K^*. The model is shown to have no difficulty in reproducing the polarization observed in pion and kaon dissociation to any final state. Other properties of the model are discussed.

1. *Introduction*

The pomeron current model is a unified phenomenological description of diffractive interactions in which pomeron exchange is represented by an effective current-current interaction. The pomeron is treated as a particle with spin one and large mass which couples to a nonconserved current V_μ. This pomeron current V_μ determines the properties of diffractive amplitudes in the same way that the electromagnetic current determines the properties of photon exchange amplitudes. For this model to be successful it is necessary that (a) diffractive interactions factorize, and (b) the approximations inherent in treating the pomeron as a spin-one object are not too drastic. Because of the small slope of the Regge pomeron trajectory this latter approximation seems to be a rather good one for limited momentum transfer (say $|t| < 1\ \mathrm{GeV}^2$), but not for large momentum transfer. Thus the model only purports to describe diffractive interactions in the region of limited $|t|$. All types of diffractive reaction are included in this description, however.

This type of model made its first appearance in the paper by Wu and Yang [1], who showed that the shape of the pp differential cross section can be fairly well represented

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by the fourth power of the nucleon electromagnetic form factor. They made no attempt to take spin effects into account; the pp elastic amplitude was simply written as a product of form factors with a fixed s dependence. Nevertheless, the essential ingredients are all present in the Wu-Yang model. It is easy to include spin effects by reinterpreting the Wu-Yang model as a current-current model and then giving the external particles spin. For nonzero spin the matrix elements of the effective current operator contain several form factors, and complicated spin-dependent scattering amplitudes can be constructed. The current--current form of the Wu-Yang model has been discussed by many authors [2–6], generally under the assumption that V_μ is a conserved current.

This assumption is motivated by an analogy with photon exchange (the so-called pomeron-photon analogy) which has a certain appeal. It is attractive to think of hadrons as extended objects which are probed by spin-one particles, the photon and the pomeron respectively, in electron-hadron scattering and diffractive hadron-hadron scattering. If this analogy makes any sense then it is possible that the pomeron current operator V_μ resembeles the $I = 0$ electromagnetic current operator.

Further progress along these lines was discouraged by the observation [4] that if V_μ in conserved then all *inelastic* diffractive cross sections must vanish in the forward direction. This contradicts experiment, and obviously the model with a conserved current must be augmented in some way. The simplest extension of the model which preserves factorization is to include in the pomeron current V_μ a second, nonconserved component. Then inelastic diffractive reactions receive a contribution in the forward direction, and the main difficulty with the older model is removed. This version of the Wu-Yang model, with each diffractive vertex having two components rather than one, was introduced by the present author [7]. A qualitative formulation of the same model in terms of valence and sea quark interactions at the diffractive vertex was suggested at the same time [7]. We turn now to a description of this model.

(a) Current-current formulation

A diffractive reaction $a+b \to c+d$ (here the final states c, d may be partial-wave states of multiparticle systems) is described by the s-channel amplitude

$$f^{(s)}_{cdab} = \lambda\langle c|V_\mu|a\rangle \langle d|V^\mu|b\rangle, \tag{1.1}$$

where λ is a constant. The pomeron current V_μ has two components

$$V_\mu = V_{1\mu}+V_{2\mu}, \tag{1.2}$$

where $V_{1\mu}$ is conserved and $V_{2\mu}$ is nonconserved. *Inelastic* matrix elements of $V_{1\mu}$ vanish in the forward direction, but forward $V_{2\mu}$ matrix elements do not vanish (unless parity invariance forces them to do so). However, both $V_{1\mu}$ and $V_{2\mu}$ contribute to *elastic* vertices in the forward direction. A very important assumption we shall make is that $V_{1\mu}$ corresponds to a central distribution of matter at the diffractive vertex, while $V_{2\mu}$ corresponds to a peripheral matter distribution. This assumption is most directly stated in terms of the form factors associated with the vertex. The central current $V_{1\mu}$ has form factors which resemble electromagnetic form factors (that is, roughly exponential, and not falling extremely rapidly with t). The peripheral curent $V_{2\mu}$ has form factors with a much sharper t dependence. This model shares certain features with the so-called two-component models of the pomeron [8]. However, it goes beyond these models by providing a direct connection between elastic and inelastic diffractive scattering.

By introducing an isoscalar nonconserved current $V_{2\mu}$ into the model we are implicitly extending the pomeron-photon analogy. The new current might be related to the $I = 0$ component of the weak neutral current. However, we do not wish to pursue this speculation here.

(b) Quark-parton formulation

We shall use the following quark-parton picture of hadrons to gain insight into the physical meaning of the operator current model just described. A hadron consists of valence quarks and a sea of $q\bar{q}$ pairs. The sea quarks correspond to a restoring force holding the valence quarks together, and we assume these sea quarks are centrally distributed. But the valence quarks, which would not be confined without this restoring force, are assumed to be peripherally distributed. This explains why πN charge exchange and other reactions which involve valence quark exchange are strongly peripheral. (Our assumption may seem to conflict with the standard quark model, where the lowest-energy wave functions tend to concentrate the (valence) quarks toward the center of a hadron. However, the quark parton model is a scattering model and not a bound-state model, and a hadron at rest has no meaning in the quark-parton description. Conventional quark model results certainly cannot be translated without change into the language we are using here). Thus, by assumption, sea quarks and valence quarks correspond to central and peripheral distributions of matter within a hadron. We identify the diffractive interaction of sea quarks and valence quarks at a diffractive vertex with the currents $V_{1\mu}$ and $V_{2\mu}$ respectively.

Although this identification does not have a clear theoretical motivation it is by no means arbitrary, as we see from the following argument. At an inelastic vertex $V_{1\mu}$ represents the interaction of a sea $q\bar{q}$ pair with a pomeron which breaks this pair apart. The two ex-sea quarks combine with the valence quarks already present to form two hadrons. But this process cannot occur for vanishing t, because a massless pomeron has zero four momentum and hence it cannot break apart the $q\bar{q}$ pair. This corresponds to $V_{1\mu}$ being conserved, so that forward inelastic $V_{1\mu}$ matrix elements vanish. The identification of $V_{1\mu}$ with a sea quark interaction is therefore quite plausible. Valence quarks, of course, scatter as independent objects and they can be scattered even by pomerons with very small four momentum. This corresponds to $V_{2\mu}$ being nonconserved. Note that for elastic vertices, $V_{1\mu}$ contributes in the forward direction because the $q\bar{q}$ pairs are not broken apart in this case and there is no reason why this interaction cannot occur at $t = 0$.

In Refs [7] we have discussed the qualitative agreement of our model with experiment. Referring the reader to these papers for commentary on elastic scattering and nucleon dissociation, we shall briefly describe here the pion and kaon dissociation data which our model is supposed to explain. (See Refs [9–11] for data on $\pi \to (3\pi)$, $(K\bar{K}\pi)$ and Refs [11–17] for data on $K \to (\pi\pi K)$, $(K\bar{K}K)$). Generally speaking, the pion and kaon data are quite similar.

(1) Diffractive transitions which do not change naturality dominate the pion and kaon dissociation data. However, final states with natural parity are also present, and their subordinate role can be partly, if not entirely, attributed to the fact that π and K dissociation into natural parity cannot proceed in the forward direction (because of factorization and parity invariance [18, 19]).

(2) There is a strong dependence of the slope of the near-forward cross section on the mass M of the final (3π), $(K\bar{K}\pi)$, $(\pi\pi K)$ or $(K\bar{K}K)$ system. In the $\pi \to (3\pi)$ case the slope is $B \approx 12$–$14\ \mathrm{GeV}^{-2}$ near the effective mass threshold, while for masses several hundred MeV larger $B \approx 5\ \mathrm{GeV}^{-2}$. Near mass threshold the cross section is approximately exponential for $|t| < 0.4\ \mathrm{GeV}^2$; for larger $|t|$ the cross section is flatter. The forward peak (with $|t| < 0.4\ \mathrm{GeV}^2$) is very sensitive to the final mass M, and it decreases in height and flattens as M increases. The cross section for $|t| > 0.4\ \mathrm{GeV}^2$ is much less sensitive to M, and when M has reached the mass range where the forward slope stops decreasing, the entire cross section out to the largest $|t|$ values is roughly exponential.

The pomeron current model offers a qualitative explanation of this slope-mass effect [7]. Matrix elements of $V_{2\mu}$ (which dominate the small-t region) are necessarily quite sensitive to the final mass, and they decrease rapidly as this mass increases. Matrix elements of $V_{1\mu}$ (which dominate the region with larger $|t|$) are not very sensitive to the final mass. This explanation, which is made possible by the two-component nature of the diffractive vertex, has yet to be tested by detailed numerical fits to data. But it is improbable that one cannot achieve reasonable fits in this way.

(3) t-channel helicity conservation (TCHC) is observed in $\pi \to \pi^*$. There is violation at the ten percent level with one unit of helicity flip allowed. Two or more units of t-channel helicity flip are not observed in any partial wave.

For any transition $\pi \to \pi^*$ or $K \to K^*$ to a state with unnatural parity, TCHC can easily be arranged in the pomeron current model by a suitable choice of meson form factors (see Sec. 2). It is just as easy to introduce an arbitrary amount of TCHC breaking by allowing one unit of t-channel helicity flip. Helicity flip of more than one unit in the t-channel is forbidden in this model for any value of t.

(4) For the $K \to K^*$ partial waves with unnatural parity, TCHC is also observed, with one exception. The Q region has recently been resolved into two resonances Q_1 (1300) and Q_2(1400) with the principal decay modes $K\varrho$ and $K^*\pi$ respectively [17]. SCHC is observed in $K \to Q_1$ and TCHC in $K \to Q_2$. Otherwise there seems to be no strong violation of TCHC in any partial wave.

For dissociation into 1^+ mesons it is as simple to arrange for SCHC as it is for TCHC by a choice of meson form factors (see Sec. 2). Thus the pomeron current model can accomodate Q_1 and Q_2 on the same footing. However, SCHC cannot be arranged for mesons with spin two or larger. (For such mesons SCHC is also impossible in the Regge model, of course).

(5) The natural-parity partial waves vanish in the forward direction, as they must if factorization holds. The most prominant ones are $\pi \to A_2$ and $K \to K^*(890)$, $K^*(1420)$, all of which violate the Gribov-Morrison rule (if these are truly diffractive reactions). In $\pi \to A_2$ only the helicity states $\lambda = \pm 1$ in the Gottfried-Jackson frame are populated

at the highest energies. Less accurate data indicates that the same may be true for $K \to K^*$ at sufficiently high energy. The helicity $\lambda = 0$ in the Gottfried-Jackson frame is forbidden by factorization and parity invariance [18, 19]. However, $\lambda = \pm 2$ for A_2 and $K^*(1420)$ is possible away from $t = 0$. The rather good high-energy data on $\pi \to A_2$ [9] show no indication that $\lambda = \pm 2$ helicity states become populated for nonzero t.

For natural-parity partial waves the pomeron current model allows only the Gottfried-Jackson frame helicities $\lambda = \pm 1$ (see Sec. 3). This holds for any t, in contrast to the Regge pomeron model where it is only true for $t = 0$.

There are weaknesses in the present model which one should keep in mind.
(*i*) There is no known theoretical justification or basis for a current-current description of diffractive interactions. The quark parton version of the model only slightly ameliorates this deficiency by providing some physical insight into the otherwise obscure existence of the pomeron current operator.
(*ii*) The energy dependence of the model is too inflexible because the pomeron spin is exactly equal to one. This means that all of the ln s effects observed in the present data are ignored. It is an open question whether or not these effects persist at asymptotic energies. If they do not, then the pomeron current model is a candidate for an asymptotic model.

As an illustration of the pomeron current model we shall construct amplitudes for pion and kaon dissociation into states with arbitrary spin and parity. Explicit amplitudes for $\pi \to \pi^*$ and $K \to K^*$ are given, where π^*, K^* are resonances or partial waves of multiparticle states such as 3π, $\pi\pi K$. In many respects (energy dependence, slope-mass effect, peripherality) the properties of these amplitudes are reasonable (at least, they can be arranged to be by a suitable choice of form factors). We shall concentrate on a different point, namely polarization. It will be shown that the model can easily reproduce the observed polarization in every diffractive $\pi \to \pi^*$ and $K \to K^*$ transition which has been studied experimentally.

2. *Dissociation into final states with unnatural parity*

We consider diffraction dissociation $N\pi \to N\pi^*$, $NK \to NK^*$ into final states π^*, K^* with the quantum numbers of the pion, kaon respectively. These states have arbitrary spin L and unnatural parity. The dissociation channel or s — channel is called $a+b \to c+d$ and the t-channel is called $D+b \to c+A$ where $A = \bar{a}$, $D = \bar{d}$. The diffractive transition (e. g. $\pi \to \pi^*$) is $d \to b$, and the elastic transition $N \to N$ is $a \to c$.

Let us begin with the t-channel helicity amplitudes

$$f^{(t)}_{cADb} = \lambda \langle cA|V_\mu|0\rangle \langle 0|V^\mu|Db\rangle. \tag{2.1}$$

The nucleon matrix element is

$$\langle cA|V_\mu|0\rangle = \bar{u}_c[\gamma_\mu f_N + \sigma_{\mu\nu} q^\nu q_N + q_\mu h_N]\, V_A, \tag{2.2}$$

where $q_\mu = (p_c + p_A)_{\mu_2}$ is the four momentum carried by the pomeron and $t = q^2$. In Eq. (2.2) the term $q_\mu h_N(t)$ corresponds to spin-zero exchange in the t-channel, and since this leads to amplitudes of order $O(1/s)$ relative to the dominant spin-one exchange amplitudes we shall simply drop the form factor $h_N(t)$ in the nucleon matrix element. (This does not imply that the current in Eq. (2.2) is conserved, for one can always project out a conserved component from any nonconserved current matrix element.) The meson matrix element is

$$\langle 0|V_\mu|Db\rangle = \Phi_{D\nu_1\cdots\nu_L} M_\mu^{\nu_1\cdots\nu_L}, \tag{2.3}$$

where Φ_D is the spin-L wave-function of π^* and

$$\begin{aligned} M_{\mu\nu_1\cdots\nu_L} &= p_{D\mu}(p_b\cdots p_b)_{\nu_1\cdots\nu_L} f_{\pi^*} \\ &+ g_{\mu\nu_1}(p_b\cdots p_b)_{\nu_2\cdots\nu_L} g_{\pi^*} \\ &+ q_\mu(p_b\cdots p_b)_{\nu_1\cdots\nu_L} h_{\pi^*}. \end{aligned} \tag{2.4}$$

Here the term containing $q_\mu h_{\pi^*}$ corresponds to spin-zero exchange in the t-channel, and we ignore this term.

We see that the nucleon vertex is essentially determined by two form factors $f_N(t)$, $g_N(t)$ and the meson vertex by two form factors $f_{\pi^*}(t)$, $g_{\pi^*}(t)$. In accordance with Eq. (1.2) we regard each of these form factors as a sum of two terms corresponding to central and peripheral distributions of matter at the respective vertices. The central form factors are roughly exponential, like electromagnetic form factors, while the peripheral form factors are much more steeply peaked in t and contribute only for small $|t|$. We note that the nucleon current (2.2) is effectively conserved (with h_N absent), and when calculating this matrix element it does not matter whether or not we assume V_μ is conserved. Therefore the nucleon form factors f_N, g_N are just the sums of form factors corresponding to $V_{1\mu}$ and $V_{2\mu}$ in Eq. (1.2),

$$f_N = f_{N1} + f_{N2}, \qquad g_N = g_{N1} + g_{N2}. \tag{2.5}$$

Things are more complicated for the meson form factors, however, since Eq. (2.4) corresponds to a nonconserved current V_μ. If instead V_μ were conserved we would have to make the following changes in this formula:

$$p_{D\mu} \to (t p_{D\mu} - q_\mu q\cdot p_D), \qquad q_{\mu\nu} \to (t q_{\mu\nu} - q_\mu q_\nu), \tag{2.6}$$

as well as requiring that $h_{\pi^*} = 0$. As we have already mentioned twice, terms which are proportional to q_μ correspond to spin-zero t-channel exchange, and such terms we can ignore. Thus the only effect of the changes (2.6) is to multiply the form factors in Eq. (2.4) by t. The meson form factors corresponding to the decomposition of V_μ in Eq. (1.2) are therefore

$$f_{\pi^*} \to t f_{\pi^*1} + f_{\pi^*2}, \qquad g_{\pi^*} \to t g_{\pi^*1} + g_{\pi^*2}. \tag{2.7}$$

The factors of t multiplying the $V_{1\mu}$ form factors f_{π^*1} and g_{π^*1} are the ones found by Ravndal [4]. They cause the $\pi \to \pi^*$ vertex to vanish in the forward direction if the pomeron current does not have a nonconserved component.

The calculation of the helicity amplitudes (2.1) is straight-forward. We find

$$\begin{aligned} G_D^L f_{cADO}^{(t)} &= \lambda(T_{db}/m_d\sqrt{2})^{L-1}\langle cA|V^\mu|0\rangle \\ &\times\{\delta_{0D} p_{D\mu}(T_{db}/m_d\sqrt{2}) f_{\pi^*}(t) + \sum_{n=0,\pm1} \delta_{nD}\varepsilon_{n\mu}(p_D) G_n^1 g_{\pi^*}(t)\}, \end{aligned} \tag{2.8}$$

where

$$G^L_m = [(2L)!/(L+m)!(L-m)!]^{1/2}, \tag{2.9}$$

$$T^2_{db} = t^2-2t(m^2_d+m^2_b)+(m^2_d-m^2_b)^2. \tag{2.10}$$

Eqs (A1), (A2) in Appendix A lead immediately to the formula (2.8). Note that t-channel helicity is conserved if $g_{\pi*} = 0$; otherwise t-channel helicity can flip by one unit. These statements hold for arbitrary final spin L.

Inserting the matrix element (2.2) into Eq. (2.8) and keeping only the leading order terms in s we find (here m is the nucleon mass)

$$G^L_1 f^{(t)}_{cc\pm 10} \approx \pm(-)^{1/2-c}\sqrt{2}\, s\sqrt{|t|}\,\lambda(t-4m^2)^{-1/2}$$
$$\times(T_{db}/m_d\sqrt{2})^{L-1}(1/T_{db})g_{\pi*}[2m(f_N-2mg_N)+(4m^2-t)g_N], \tag{2.11}$$

$$G^L_0 f^{(t)}_{cc00} \approx (-)^{1/2+c}s\lambda(T_{db}/m_d\sqrt{2})^{L-1}(1/m_d\sqrt{2}\,T_{db})\,(t-4m^2)^{-1/2}$$
$$\times[T^2_{db}f_{\pi*}+2(t+m^2_d-m^2_b)g_{\pi*}]\,[2m(f_N-2mg_N)+(4m^2-t)g_N], \tag{2.12}$$

$$G^L_1 f^{(t)}_{c-c\pm 10} \approx \mp\sqrt{2}\, s|t|\lambda(t-4m^2)^{-1/2}(1/T_{db})$$
$$\times(T_{db}/m_d\sqrt{2})^{L-1}g_{\pi*}(f_N-2mg_N), \tag{2.13}$$

$$G^L_0 f^{(t)}_{c-c00} \approx s\sqrt{|t|}\,\lambda(T_{db}/m_d\sqrt{2})^{L-1}(1/m_d\sqrt{2}\,T_{db})$$
$$\times(t-4m^2)^{-1/2}(f_N-2mg_N)\,[T^2_{db}f_{\pi*}+2(t+m^2_d-m^2_b)g_{\pi*}]. \tag{2.14}$$

The cross section is

$$\frac{d\sigma}{dt} \approx \frac{1}{16\pi}\frac{|\lambda|^2}{m^2_d(2L)!}\frac{1}{T^2_{db}}\left(\frac{T_{db}}{m_d\sqrt{2}}\right)^{2L-2}(|f_N|^2+|t|\,|g_N|^2)$$
$$\times\{(L!)^2|T^2_{db}f_{\pi*}+2(t+m^2_d-m^2_b)g_\pi|^2+8|t|m^2_d(L+1)!(L-1)!|g_{\pi*}|^2\}. \tag{2.15}$$

From Eq. (1.1) we calculate the s-channel helicity amplitudes

$$G^L_d f^{(s)}_{cda0} = \lambda(T_{db}/m_d\sqrt{2})^{L-1}(-)^d\langle c|V^\mu|a\rangle\,\{-p_{d\mu}L![(L+d)!(L-d)!]^{-1/2}$$
$$\times(T_{db}/m_d\sqrt{2})d^L_{0d}(\chi_d)f_{\pi*}+\sum_{n=0,\pm 1}\varepsilon_{n\mu}(p_d)\,(-)^n G^1_n(L-1)!$$
$$\times[(L-1+d-n)!(L-1-d+n)!]^{-1/2}d^{L-1}_{0,d-n}(\chi_d)g_{\pi*}\}, \tag{2.16}$$

where Eqs (A2) and (A3) have been used and the angle χ_d is defined by

$$S_{cd}T_{db}\sin\chi_d = 2m_d\sqrt{\phi}\,, \tag{2.17}$$

$$S_{cd}T_{db}\cos\chi_d = -(s-m^2+m^2_d)\,(t+m^2_d-m^2_b)+2m^2_d(m^2_d-m^2_b), \tag{2.18}$$

$$S^2_{cd} = s^2-2s(m^2+m^2_d)+(m^2-m^2_d)^2. \tag{2.19}$$

We are particularly interested in helicity amplitudes for $L = 1$. For large s these are

$$f^{(s)}_{c\pm 1,c0} \approx \pm \frac{\lambda}{\sqrt{2}} s\sqrt{|t|}\, f_{\pi^*} f_{\mathrm{N}}, \tag{2.20}$$

$$f^{(s)}_{c\pm 1,-c0} \approx \pm(-)^{1/2-c} \frac{\lambda}{\sqrt{2}} s|t|\, f_{\pi^*} g_{\mathrm{N}}, \tag{2.21}$$

$$f^{(s)}_{c0,c0} \approx \frac{\lambda s}{2m_d} f_{\mathrm{N}}[(t+m_d^2-m_b^2)f_{\pi^*}+2g_{\pi^*}] \tag{2.22}$$

$$f^{(s)}_{c0,-c0} \approx (-)^{1/2-c} \frac{\lambda s}{2m_d} \sqrt{|t|}\, g_{\mathrm{N}}[(t+m_d^2-m_b^2)f_{\pi^*}+2g_{\pi^*}]. \tag{2.23}$$

Note that s-channel helicity is conserved at the nucleon vertex if $g_{\mathrm{N}} = 0$. (If V_μ were the $I = 0$ electromagnetic current then g_{N} would be proportional to the nucleon isoscalar anomalous magnetic moment, which is small). At the meson vertex SCHC holds if $f_{\pi^*} = 0$. Therefore, in the case $L = 1$ we can arrange for SCHC as easily as we can TCHC. But this is only true for $L = 1$.

3. *Dissociation into final states with natural parity*

Now we turn to diffraction dissociation into final states π^*, K^* with natural parity and arbitrary spin L.

The t-channel helicity amplitudes are given by Eq (2.1) with the nucleon vertex as in Eq. (2.2) and the meson vertex as in Eq. (2.3), but now

$$M_{\mu\nu_1\dots\nu_L} = \varepsilon_{\mu\nu_1\alpha\beta} p_D^\alpha p_b^\beta (p_b \dots p_b)_{\nu_2\dots\nu_L} f_{\pi^*}(t). \tag{3.1}$$

There is only one form factor at this vertex. The tensor (3.1) has the same form for conserved and nonconserved currents. Therefore the form factor which corresponds to the current decomposition (1.2) is just a sum

$$f_{\pi^*} = f_{\pi^* 1} + f_{\pi^* 2} \tag{3.2}$$

of central and peripheral form factors.

Calculating the t-channel helicity amplitudes we find

$$G_D^L f^{(t)}_{cA\pm 10} = \pm i \frac{\lambda}{\sqrt{2}} m_d (T_{db}/m_d\sqrt{2})^L \varepsilon_{\pm 1\mu}(p_D) \langle cA|V^\mu|0\rangle f_{\pi^*}, \tag{3.3}$$

while $f^{(t)}_{cAD0} = 0$ for $D \neq \pm 1$.

Here we have used Eqs (A1), (A2), (A4). In the current-current model only the t-channel π^* helicity states $D = \pm 1$ are populated. t-channel helicity of two or more units is strictly forbidden in this model, and the helicity state $D = 0$ is forbidden by parity invariance [18, 19]. This leaves only the $D = \pm 1$ helicity states. Experimentally these are the only ones observed in A_2 production at 40 GeV [9].

In the Regge model one obtains the same result at $t = 0$. However, away from $t = 0$ the helicity states $|D| > 1$ can also play a role because the Regge pomeron spin is not exactly one. A good test of the Regge model would be an experiment on A_2 polarization for $|t| > 0.5$ GeV^2 (the present data at 40 GeV [9] is limited to the range $0.17 < |t| < 0.33$ GeV^2). If the helicity states $D = \pm 2$ in the Gottfried–Jackson frame become populated for large $|t|$, then the changing spin of the Regge pomeron plays an essential role. If not, then the pomeron current model should provide an adequate description for $|t| > 0.5$ GeV^2 as it seems to do for smaller $|t|$ values. In this same experiment one would also test factorization at large $|t|$ by measuring the contribution of the $D = 0$ state in the Gottfried–Jackson frame.

Keeping only the leading terms in s we obtain explicit formulas for the helicity amplitudes;

$$G_1^L f_{cc\pm 10}^{(t)}$$

$$\approx i(-)^{1/2+c} \frac{\lambda s\sqrt{|t|}}{\sqrt{2}\,(t-4m^2)^{1/2}} \left(\frac{T_{db}}{m_d\sqrt{2}}\right)^{L-1} f_{\pi^*}[2m(f_N - 2mg_N) + (4m^2-t)g_N], \tag{3.4}$$

$$G_1^L f_{c-c\pm 10}^{(t)} \approx i \frac{\lambda s|t|}{\sqrt{2}\,(t-4m^2)^{1/2}} \left(\frac{T_{db}}{m_d\sqrt{2}}\right)^{L-1} f_{\pi^*}(f_N - 2mg_N). \tag{3.5}$$

The cross section is

$$\frac{d\sigma}{dt} \approx \frac{|t|}{8\pi} \frac{(L+1)!(L-1)!}{(2L)!} \left(\frac{T_{db}}{m_d\sqrt{2}}\right)^{2L-2} |\lambda f_{\pi^*}|^2[|f_N|^2 + |t|\,|g_N|^2]. \tag{3.6}$$

In Ref. [11] the A_2 production cross section with the factor $|t|$ removed is compared with the ones for A_1 and A_3 production (these do not vanish in the forward direction). Little difference is found: all three are roughly exponential, with the A_2 exponent lying between the larger A_1 exponent and the smaller A_3 exponent, in accordance with the usual slope-mass effect. Evidently the production mechanisms for A_1, A_2 and A_3 are quite similar. In the pomeron current model this tells us that the form factor f_{A_2} for $\pi \to A_2$ is similar to the ones for $\pi \to A_1, A_3$ (here we mean f_{π^*} since TCHC implies $g_{\pi^*} \approx 0$ in A_1 A_3 production). In particular, the corresponding $V_{2\mu}$ form factors $f_{\pi^* 2}$ must be similar because the $V_{1\mu}$ form factors do not contribute in $\pi \to A_{1,3}$ for small $|t|$ (see Eq. (2.7)).

APPENDIX A

Some formulas needed to obtain results given in the text are the following:

$$(p_b \cdots p_b)^{\nu_1 \cdots \nu_L} \Phi_{D\nu_1 \cdots \nu_L}(p_D) = \delta_{0D}(T_{db}/m_d\sqrt{2})^L L!/\sqrt{(2L)!}\,, \tag{A1}$$

$$G_D^L \Phi_{D\nu_1 \cdots \nu_L}(p_D) = \sum_{n=0,\pm 1} G_n^1 G_{D-n}^{L-1} \varepsilon_{n\nu_1}(p_D) \Phi_{D-n,\nu_2 \cdots \nu_L}(p_D), \tag{A2}$$

$$(p_b \cdots p_b)^{\nu_1 \cdots \nu_L} \Phi_{d\nu_1 \cdots \nu_L}(p_d) = (-)^d d_{0d}^L(\chi_d)\,(T_{db}/m_d\sqrt{2})^L L!/\sqrt{(2L)!}\,, \tag{A3}$$

$$\varepsilon^{\mu\nu\alpha\beta} \varepsilon_{n\nu}(p_D) p_{D\alpha} p_{b\beta} = \pm\delta_{\pm 1,n} \frac{i}{2} T_{db} \varepsilon_{\pm 1}^{\mu}(p_D). \tag{A4}$$

APPENDIX B

Here we give a general formula for the construction of t-channel helicity amplitudes in the pomeron current model for any two-body diffractive process $ab \to cd$

$$f_{cADb}^{(t)} = -\lambda d_{D-b,c-A}^{1}(\theta_t) g_{cA}(t) g_{Db}(t)$$

where for large s

$$\cos\theta_t \approx i \sin\theta_t \approx -2s|t|/T_{ca}T_{db}.$$

The form factors $g_{cA}(t)$ and $g_{Db}(t)$ which determine the vertices $a \to c$ and $b \to d$ are defined as follows.

$$T_{cA}^{m} = \varepsilon_{m\mu}(q) \langle cA|V^{\mu}|0\rangle,$$

$$T_{Db}^{n} = \varepsilon_{n\mu}^{*}(q) \langle 0|V^{\mu}|Db\rangle.$$

Then introduce two different t-channel CM frames Q_{cA} and Q_{Db} (these are related by a rotation about the y-axis through an angle $-\theta_t$);

$$Q_{cA}: \quad \vec{q} = \vec{p}_c + \vec{p}_A = 0, \quad \hat{p}_c = -\hat{p}_A = \hat{q} = \hat{z}, \quad T_{cA}^{m} = \delta_{m,c-A} g_{cA}(t);$$

$$Q_{Db}: \quad \vec{q} = \vec{p}_D + \vec{p}_b = 0, \quad \hat{p}_D = -\hat{p}_b = \hat{q} = \hat{z}, \quad T_{Db}^{n} = \delta_{n,D-b} g_{Db}(t).$$

Here carats denote unit three-vectors, and $\hat{q}$ is the direction chosen for the spin quantization axis of the polarization vectors $\varepsilon_{n\mu}(q)$ associated with pomeron exchange. From angular momentum conservation the functions T_{cA}^{m} and T_{Db}^{n} have the simple forms shown above in the frames Q_{cA} and Q_{Db}, respectively. This defines the form factors $g_{cA}(t)$ and $g_{Db}(t)$ as linear combinations of the covariant form factors which determine the matrix elements $\langle cA|V_{\mu}|0\rangle$ and $\langle 0|V_{\mu}|Db\rangle$.

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