PHYSICAL REVIEW D VOLUME 22, NUMBER 7 1 OCTOBER 1980

# Diffractive dissociation processes in a field-theoretic quark model

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(Received 19 November 1979)

We consider here diffractive dissociation processes like $A+B\rightarrow A+X$ in a field-theoretic quark model with a phenomenological current-current interaction in quark space used earlier for discussing diffraction scattering in the same model. The process $A+B\rightarrow A+X$ is interpreted as being diffractive for the hadron $A$ with correlation to the form factor of $A$, and is incoherent for the hadron $B$ with correlation to the structure function of $B$. These assumptions in the *quark model* yield that the cross sections $d\sigma/dt\,dM^2$ have the same structure as is otherwise derived on the basis of the more established triple-Pomeron coupling, which could probably be partly anticipated since the Pomeron corresponds to a spin-1 object with vacuum quantum numbers. The results are compared with experiments, and broad agreement at high energies is observed. With these ideas, in the region of small mass of the dissociated system, the contributions to the cross section from diffractive dissociation and from diffractive production of resonances can be separated, where a limitation for the incoherent scattering in quark-parton models is used.

## I. INTRODUCTION

We shall consider here diffractive dissociation processes[1] in a field-theoretic quark model proposed recently.[2] The model has been considered for many coherent hadronic processes with some success.[3,4] Here in particular a phenomenological current-current form of interaction in quarks space appears to be useful for obtaining diffraction scattering of hadrons in a universal manner with the diffraction slopes related to the size of the hadronic wave functions as determined from their static properties[2-4] and in rough agreement with experiments.[4] For this reason we examine here diffractive dissociation processes for scattering of hadrons with the same interaction.

Diffractive dissociation of hadrons was predicted[1] a long time back in analogy with diffractive scattering in optics. Also, an impact-parameter picture[5] or duality-diagram calculations[6] have been used for understanding them. However, the accepted explanation for these diffractive processes is triple-Regge phenomenology,[7] also applicable to the more general processes like $A+B\rightarrow C+X$. In the present analysis, we examine the diffractive dissociation processes $A+B\rightarrow A+X$ only, governed at high energies by the triple-Pomeron coupling.[8] This appears to be reasonable since the earlier analysis for diffraction scattering[4] corresponds to the Pomeron exchange.

We note that a current-current form of interaction in quark or in hadronic space has been earlier utilized by many authors[9] for diffraction scattering, diffractive photoproduction, or even for large-angle scattering of hadrons. Our approach here will follow, however, mostly that of Ravndal.[10]

In Sec. II we calculate the cross-sections for diffractive dissociation processes in the quark-parton model, and obtain the same theoretical form for them as in Regge-pole phenomenology with the triple-Pomeron contribution.[8] This in itself appears interesting since the two methods are quite different. In Sec. III we consider possible limitations of quark-parton ideas in the context of incoherent hadronic processes on some physical grounds and discuss some of the implications of the present ideas.

## II. DIFFRACTIVE DISSOCIATION OF HADRONS

We shall now consider the diffractive dissociation processes like $A+B\rightarrow A+X$ with the phenomenological interaction[4]

$$\mathcal{H}_I(x)=f_V\, J_s^{\mu}(x)\, J_{s\mu}(x)\,, \qquad (2.1)$$

where, unlike Ref. 4, we shall not use the explicit form of the strong current $J_s^{\mu}(x)$ or the hadronic states except for gross features like quark additivity and form factors. We had taken earlier[4,2]

$$J_s^{\mu}(x)=\overline{Q}^g_{\alpha i}(x)\,\gamma^{\mu}\,Q^g_{\alpha i}(x)-\bar{\tilde{Q}}^g_{\alpha i}(x)\,\gamma^{\mu}\,\tilde{Q}^g_{\alpha i}(x)\,, \qquad (2.2)$$

with summation over the flavor and color indices $\alpha$ and $i$. Such conclusions as we state here can be easily verified as earlier[2,4] with the appropriate hadronic states and the quark field operators.

In order to fix our ideas, we shall consider the diffractive dissociation process corresponding to $p+p\rightarrow p+X$ in some detail and state the results in the cases of the other similar processes.

### A. $p+p\rightarrow p+X$

With $p_1r_1$, $p_2r_2$, and $p_2'r_2'$ as the momenta and spins of the initial protons and the final proton, and with the interaction (2.1), we associate the matrix element for the diffractive dissociation process as[2]

$$M_{fi} = -i(2\pi)^4 2f_V \langle n|J_s^\mu(0)|p_1 r_1\rangle \times \langle p_2' r_2'|J_{s\mu}(0)|p_2 r_2\rangle . \qquad (2.3)$$

In the above, $n$ is a specific "dissociation" channel of the proton. We may regard $n$, e.g., as a system of three free quarks,[11] or three free quarks and a gluon,[12] or as a more complicated system which finally hadronizes forming the observed hadrons. This corresponds to incoherent scattering in quark space.[13] The second matrix element on the right-hand side of (2.3) on the other hand corresponds to a coherent hadronic process.[4,9,10] Thus we have assumed that with the interaction (2.1) the diffractive dissociation process is partly incoherent[11,13] and is partly coherent, and is given in a quantitative manner by (2.3). If we have resonance production, $n$ in (2.3) will correspond to a resonance. This will be excluded in the present discussions for diffractive dissociation processes.

Now, the cross section given by (2.3) becomes

$$\sigma(p_1,p_2\to X,p_2') = \frac{4\pi^2}{V_{\rm rel}} \int\sum_n (2\pi)^8\, 4f_V^2\, d^3p_2'\, \delta^4(p_1+p_2-p_2'-p_n)\, \langle n|J_s^\mu(0)|p_1 r_1\rangle \tfrac{1}{2}\langle p_1 r_1|J_s^\nu(0)|n\rangle \times \langle p_2' r_2'|J_{s\mu}(0)|p_2 r_2\rangle \tfrac{1}{2}\langle p_2 r_2|J_{s\nu}(0)|p_2' r_2'\rangle . \qquad (2.4)$$

For the coherent part of the matrix element in (2.4), we substitute with appropriate spin summations

$$\tfrac{1}{2}\langle p_2' r_2'|J_{s\mu}(0)|p_2 r_2\rangle\langle p_2 r_2|J_{s\nu}(0)|p_2' r_2'\rangle = 9\times(2\pi)^{-6}\frac{1}{p_2^0 p_2'^0} H_{p\mu\nu}(p_2,p_2') , \qquad (2.5)$$

where[14] for the proton we have after some simplifications

$$H_{p\mu\nu}(p_2,p_2') = \left(1-\frac{t}{4m^2}\right)^{-1}(G_E^2 - t\,G_M^2)\,\tfrac{1}{4}(p_{2\mu}+p_{2\mu}')(p_{2\nu}+p_{2\nu}') + m^2 G_M^2 (g_{\mu\nu}t - q_{2\mu}q_{2\nu}) . \qquad (2.6)$$

In (2.6), $G_E(t)$ and $G_M(t)$ are the electric and magnetic form factors and we have taken these form factors for the strong current to be the same as those for the electromagnetic current.[4] The factor 9 in (2.5) may be noted and is due to the fact that the proton has three quarks[2,4,10] and is equivalent to the earlier quark counting rule of quark model for hadronic reactions.

For the incoherent part in diffractive dissociation we substitute[15]

$$\tfrac{1}{2}\sum_n \delta(p_1+q_2-p_n)\,\langle n|J_s^\mu(0)|p_1 r_1\rangle\langle p_1 r_1|J_s^\nu(0)|n\rangle = 3\times(2\pi)^{-6}\frac{m}{p_1^0} W_s^{\mu\nu}(p_1,q_2) , \qquad (2.7)$$

where $W_s^{\mu\nu}$ is the structure function corresponding to the strong current. Since the process is incoherent, here we have a factor 3 on the right-hand side of (2.7) instead of 9 corresponding to the three quarks of the single proton.[11]

With (2.5) and (2.7), we obtain from (2.4) that

$$\sigma(p_1,p_2\to X,p_2') = \frac{4\pi^2}{V_{\rm rel}} \int (2\pi)^{-4}\, 4f_V^2 \frac{3m}{p_1^0} W_s^{\mu\nu}(p_1,q_2)\frac{9}{p_2^0 p_2'^0} H_{p\mu\nu}(p_2,p_2')\, d^3p_2' . \qquad (2.8)$$

With a conventional identification[15] of $W_s^{\mu\nu}(p_1,q_2)$ and using (2.6) we get, with $F_{2p}$ and $F_{1p}$ as the structure functions of the proton,

$$W_s^{\mu\nu}(p_1,q_2) H_{p\mu\nu}(p_2,p_2') = \frac{F_{2p}}{m^2\nu}\left[\left(1-\frac{t}{4m^2}\right)^{-1}(G_E^2 - tG_M^2)\,\tfrac{1}{4}(s-m\nu-2m^2)^2 - m^4 G_M^2(\nu^2-t)\right] - mF_{1p}(G_E^2 + 2tG_M^2) . \qquad (2.9)$$

Further, choosing $\vec{p}_2 = 0$, we get

$$\frac{1}{V_{\rm rel}}\frac{d^3p_2'}{p_1^0 p_2^0 p_2'^0} \equiv \frac{2\pi}{s^2}\, dt\, dm^2 , \qquad (2.10)$$

where

$$M^2 = (p_1+q_2)^2 = m^2 + 2m\nu + t \qquad (2.11)$$

is the square of the mass of the "dissociated" system of hadrons. We now see that the contribution is suppressed like a square of the form factor in (2.9). We make an exponential approximation for this for small $t$ corresponding to the form factor[10] for $p$-$p$ elastic scattering, and consider the approximation that $M^2 \ll s$. We then obtain from (2.4) that

$$\frac{d\sigma}{dt\,dM^2}(pp \to Xp) = \frac{27 f_V^2}{\pi}\frac{F_{2p}}{2m\nu}\left(1-\frac{2m\nu}{s}\right)\exp(\tfrac{1}{3}R_p^2 t)\,, \quad (2.12)$$

where with $R_p$ as the charge radius of the proton we have approximated (2.9) by

$$W_s^{\mu\nu}(p_1, q_2)\, H_{p\mu\nu}(p_2, p_2') \simeq s^2 \frac{F_{2p}}{4m^2\nu}\left(1-\frac{2m\nu}{s}\right)\exp(\tfrac{1}{3}R_p^2 t). \quad (2.13)$$

If we substitute the Feynman variable $x_F$ as

$$x_F = 1-(M^2/s)\,, \quad (2.14)$$

(2.12) then becomes

$$\frac{d\sigma}{dt\,dM^2}(pp \to Xp) \simeq \frac{27 f_V^2}{\pi}\frac{F_{2p}}{M^2-m^2}\, x_F \exp(\tfrac{1}{3}R_p^2 t), \quad (2.15)$$

which has the same form as for the triple-Pomeron vertex contribution when $M^2 \gg m^2$, and we remember that we should have $F_{2p}(x) = F_{2p}(0)$ as constant with the Bjorken variable $x$ in the present limit being almost zero. We note that with (2.2) in mind the contribution from the strong current to the corresponding structure function need not be the same as that of the electromagnetic structure function, except probably the valence contribution.[11] Since the contribution in (2.15) from the sea[12] we expect the constant $F_{2p}(0)$ to be different from that of the limiting value of the structure function for small $x$, as obtained from deep-inelastic lepton-hadron scattering.

We now note that for the diffractive dissociation process in (2.15), the differential cross section has been parametrized from experimental data as[16]

$$\frac{d\sigma}{dt\,dM^2} = \frac{A b_0(1+B/p_{\text{lab}})}{M^2}\exp(b_0 t) \quad (2.16)$$

with $Ab_0 \simeq 3.5 \pm 0.2$ $mb/\text{GeV}^2$, $b_0 \simeq 6.5 \pm 0.3$ $\text{GeV}^{-2}$, and $B \simeq 54 \pm 16$ GeV. As a comparison, (2.15) yields that[17] $b_0 = \frac{1}{3}R_p^2 \simeq 6$ $\text{GeV}^{-2}$. The constant $B$ in (2.16) is an effect of nonleading-order contributions not retained here, and its contribution vanishes in the high-energy limit. $F_{2p}$, which was regarded as arbitrary, is now determined as

$$F_{2p} \simeq A b_0 \frac{\pi}{27 f_V^2} \simeq 0.13 \pm 0.01\,, \quad (2.17)$$

where we have used that from diffraction scattering[4] $f_V \simeq 2.8$ $\text{GeV}^{-2}$. We notice that (2.17) is a reasonable value for the low-$x$ limit of the structure function as known from lepton-hadron processes, and , as stated earlier, a difference of factor 2 is not unexpected. We consider the form factors to be the same as for electromagnetic current whereas the structure function may be different because the form factor will be given by the valence quarks,[4] whereas the structure function for small $x$ does not correspond to the valence quarks and hence may change with something like (2.2).

We thus note that the slope parameter for the $t$ dependence, the $M^2$ and $x_F$ dependence, and even the order of magnitude of $d\sigma/dt\,dM^2$ is obtained with the same assumptions as in Ref. 4 in the quark model, without any direct use of the triple-Pomeron vertex. The only parameter here adjusted is $F_{2p}$. In Fig. 1, we plot $(d\sigma/dt\,dM^2)$ $(pp \to Xp)$ in *curve* I as calculated from (2.15) against the experimental points.[16,18] We note that for intermediate $M^2$, the present description is good. For small $M^2$ the description is expected to be bad for two reasons. Firstly, from kinematic considerations the cross section must vanish when $M$ is small enough. In that case dissociation resulting from incoherent scattering with impulse approximation will be a wrong assumption and therefore (2.15) will no longer be valid. We shall discuss this aspect in the Sec. III. The other reason why (2.15) will be wrong is because, again, the assumption of dissociation will be wrong with the production of resonances in a coherent manner, which, as mentioned earlier, is not included here.

We also note that[18,19] the slope in $t$ of $d\sigma/dt\,dM^2$ for $1.5 < M^2 < 2.5$ is substantially higher than 6. This is quite consistent with the observation made above, since in this region resonances will be produced in a coherent manner parallel to the slope for diffraction scattering.[4] On the other hand, for $M^2 > 3$, the slope appears to be consistent with 6 $\text{GeV}^{-2}$ even in the recent observations[19], which demonstrates that the present naive picture with factorizable diffractive and dissociation effect as taken in the present model is quite consistent. On integrating, from (2.15) we get that

$$\int_{-0.235}^{-0.024}\frac{d\sigma}{dt\,dM^2}\,dt = \frac{27 f_V^2}{\pi}\frac{F_{2p}}{M^2}\,x_F \times \frac{3}{R_p^2}\left[\exp\left(-\frac{R_p^2}{3}\times 0.024\right) - \exp\left(-\frac{R_p^2}{3}\times 0.235\right)\right]. \quad (2.18)$$

In Fig. 2 we compare the curve as calculated from (2.18) with the experimental points from Ref. 19. As the energy increases there is progressively better agreement between the theoretical curve and the experimental points as a confirmation of the present picture.

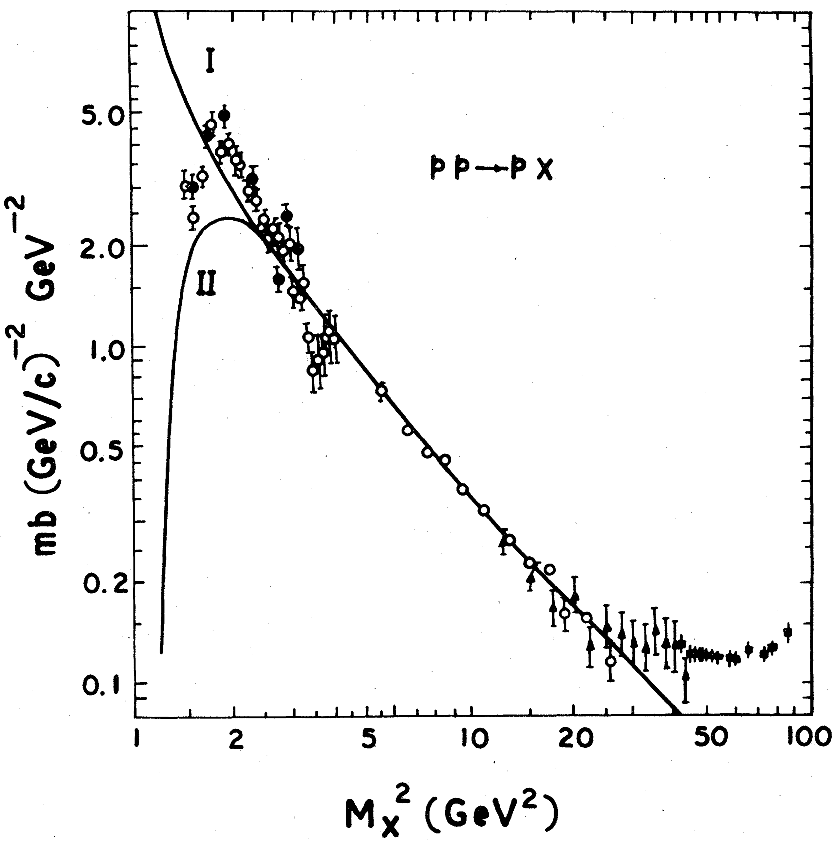


FIG. 1. $(d\sigma/dt\,dM^2)(pp\rightarrow pX)$ vs $M^2$ at small $t$ due to diffraction dissociation. Curve I is according to (2.15) and curve II is according to (3.3). The experimental points are from Ref. 16. The open circles, solid circles, solid triangles, and solid squares represent the data at laboratory energies 275, 260, 281, and 267 GeV/$c$, respectively.

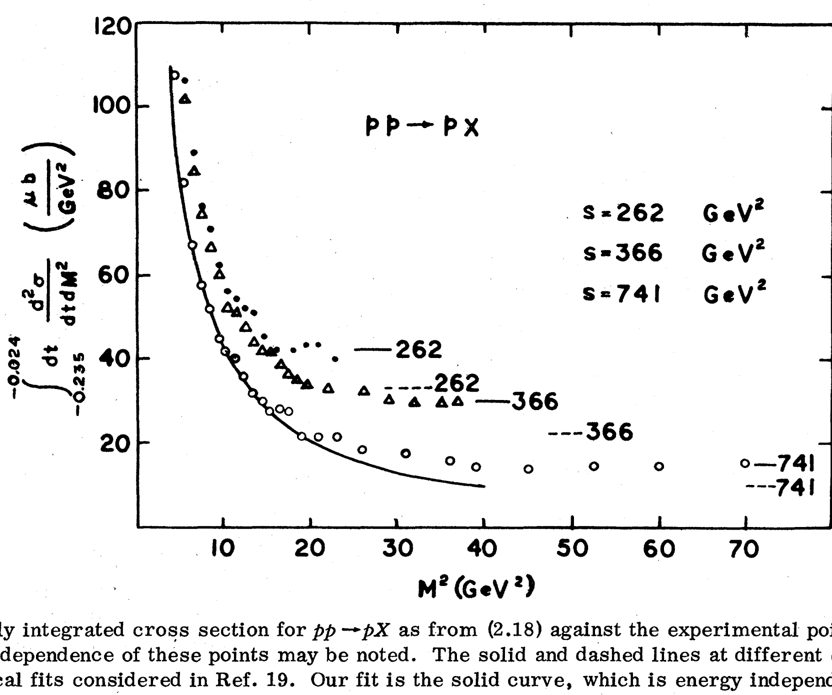


FIG. 2. Partially integrated cross section for $pp\rightarrow pX$ as from (2.18) against the experimental points of Ref. 19, where the energy dependence of these points may be noted. The solid and dashed lines at different energies correspond to phenomenological fits considered in Ref. 19. Our fit is the solid curve, which is energy independent and corresponds to the high-energy limit.

As noted in the present analysis we have only a *single* adjustable parameter $F_{2p}$, considering which the agreement with experiments appears as quite reasonable.

### B. $\pi+p\rightarrow X+p$

We next consider the diffractive dissociation process for $\pi p$ scattering when the pion dissociates. The diffractive part of the matrix element will remain unchanged, and for the dissociation of the pion, (2.7) will be replaced by

$$\sum_n \delta_4(p_1+q_2-p_n)\langle n|J_s^\mu(0)|p_1\rangle\langle p_1|J_s^\nu(0)|n\rangle = 2\times(2\pi)^{-6}\frac{m}{p_1^0}W_s^{\mu\nu}(p_1,q_2)\,. \quad (2.19)$$

We may note the change in the factor from 3 to 2 in (2.19) for the quark content of the meson with quark additivity.[4,11] As before, proceeding with the leading approximation we obtain, corresponding to (2.15),

$$\frac{d\sigma}{dt\,dM^2}(\pi p\rightarrow Xp)=\frac{18f_V^2}{\pi}\frac{F_{2\pi}}{M^2}x_F\exp(\tfrac{1}{3}R_p^2t)\,. \quad (2.20)$$

As before, in the above $F_{2\pi}$ is the only arbitrary parameter. This corresponds to the pion structure function for the limit of small Bjorken variable $x$. We thus predict that here also we shall have the same slope in $t$ as in the last subsection. Comparing with experiments[20] we find that in fact for $M^2>4$ this is so. As before, for $M^2<4$, the larger slope corresponds to the coherent production of resonances similar to the slope for diffraction scattering.[4,10] When we integrate (2.20) over $t$ we get

$$\frac{d\sigma}{dM^2}(\pi p\rightarrow Xp)=\frac{18f_V^2}{\pi}\frac{F_{2\pi}}{M^2}x_F\frac{3}{R_p^2}\,. \quad (2.21)$$

We have plotted in Figs. 3 and 4 $d\sigma/dM^2$ as calculated from (2.21) against the experimental points,[20,21] choosing $F_{2\pi}=0.18$. From other considerations of the pion structure function[22] this appears to be reasonable although we may mention again that it is really unknown and is taken as a free parameter. We may also note that from (2.21) we obtain for $4<M^2<32$ GeV$^2$ that $\sigma=1.09$ mb, whereas the experimental value quoted[20] is $0.9\pm0.2$ mb.

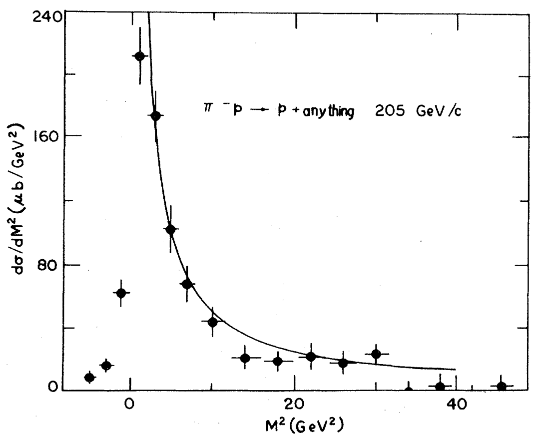


FIG. 3. $(d\sigma/dM^2)(\pi p\rightarrow Xp)$ vs $M^2$ as per (2.21) against the experimental points of Ref. 20.

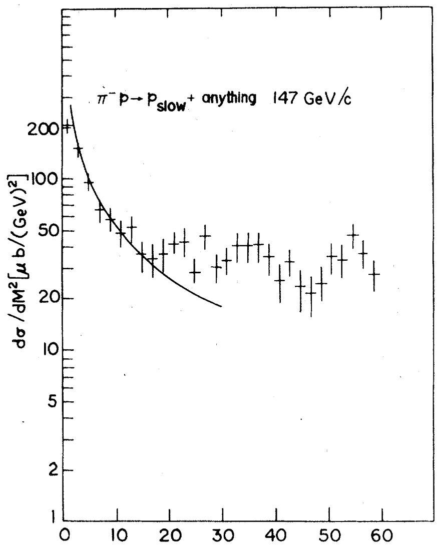


FIG. 4. $(d\sigma/dM^2)(\pi p\rightarrow Xp)$ vs $M^2$ as per (2.21) against the experimental points of Ref. 21.

### C. $\pi+p\rightarrow\pi+X$

We next consider the diffractive dissociation process for $\pi p$ scattering when the proton dissociates. As earlier, here the square of the form factor for the pion will enter the picture, along with $F_{2p}$ for the proton. With the association between the diffraction slopes and the form factors which we are using in the present crude approximations[4,10] we shall here have

$$F_\pi^2(t)=\exp(\tfrac{1}{3}R_\pi^2t)\,, \quad (2.22)$$

where $R_\pi$ is the charge radius of the pion.[23] Substituting this, as before we now obtain

$$\frac{d\sigma}{dt\,dM^2}(\pi p\rightarrow\pi X)=\frac{12f_V^2}{\pi}\frac{F_{2p}}{M^2-m^2}x_F\exp(\tfrac{1}{3}R_\pi^2t)\,, \quad (2.23)$$

which yields on integration

$$\frac{d\sigma}{dM^2}(\pi p\rightarrow\pi X)=\frac{12f_V^2}{\pi}\frac{F_{2p}}{M^2-m^2}x_F\frac{3}{R_\pi^2}\,. \quad (2.24)$$

In Fig. 5 we compare (2.24) with experimental points[21] where we take[23] $R_\pi^2=0.48$ fm$^2$. The results appear to be reasonable when we recognize, com-

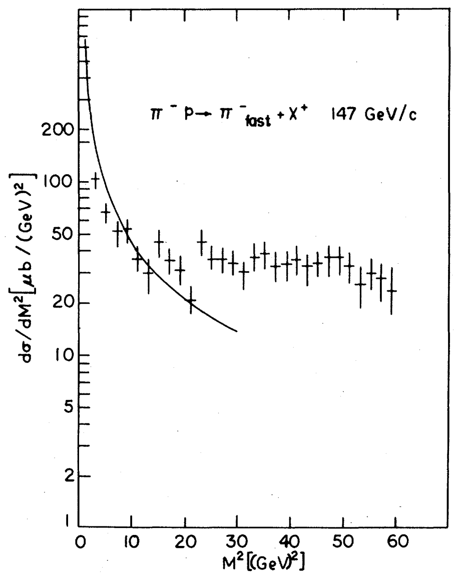


FIG. 5. $(d\sigma/dM^2)$ $(\pi p \to \pi x)$ vs $M^2$ as per (2.24) against the experimental points of Ref. 21.

paring with Fig. 2, that the energy is not yet high enough. In this subsection no constant has been adjusted.

As in the earlier subsections A and B, the results are not good for $M^2 \gtrsim 30$ GeV$^2$, which is not unexpected, since in this case in conventional models we are to go beyond triple-Pomeron coupling, and in the context of the present model we expect that also the assumption of factorizability may no longer be valid.

### D. $K+p \to X+p$

We next consider diffractive dissociation of the $K$ meson with $Kp \to Xp$ as an example of this process with SU(3)-symmetry violation. We note that with the assumption that at high energies the forward scattering amplitude is pure imaginary, we have for $\pi p$ diffraction scattering[4,10,24]

$$\frac{d\sigma}{dt}(\pi p \to \pi p)\bigg|_{t=0} = \frac{36 f_V^2}{\pi} = \frac{\sigma_t^2(\pi p)}{16\pi}, \tag{2.25}$$

whereas for $Kp$ diffraction scattering[4,10]

$$\frac{d\sigma}{dt}(Kp \to Kp)\bigg|_{t=0} = \frac{18 f_V^2(1+\theta^2)}{\pi} = \frac{\sigma_t^2(Kp)}{16\pi}. \tag{2.26}$$

We have written down (2.26), with the coupling constant $f_V$ in (2.1) changed to $\theta f_V$ when a $\lambda$ quark current is present, in order to include the effect of SU(3) symmetry violation. Now for the diffractive dissociation process with the dissociation of $K$, a fresh constant $F_{2K}$ will further enter in the same manner as $F_{2p}$ or $F_{2\pi}$ earlier. With this, parallel to (2.20) we now have

$$\frac{d\sigma}{dt\,dM^2}(Kp \to Xp) = \frac{\sigma_t^2(Kp)}{\sigma_t^2(\pi p)} \frac{18 f_V^2}{\pi} \frac{F_{2K}}{M^2} x_F \exp(\tfrac{1}{3}R_p^2 t), \tag{2.27}$$

where we have included the effect of symmetry breaking with the total cross sections, and the effect of symmetry breaking will further occur through $F_{2K}$. The overall correction of both these factors is a normalization constant.

Data are not available at *adequately* high energies[19,21,25] for the process (2.27) such that the triple-Pomeron coupling dominates.[25] However, for small $M^2/s$ ($\lesssim 0.2$) the qualitative ($1/M^2$) behavior predicted by (2.27) appears to be there, and further the slope in $t$ for this region is also about[25] 6 GeV$^{-2}$ as predicted. It is clear in the analysis of Ref. 25 that to fit the data at these energies one must include many parameters with many Regge exchanges, which in the context of the present model corresponds to more complicated reactions not included in the phenomenological Hamiltonian taken in (2.1) corresponding to the Pomeron coupling. Thus, in the context of Fig. 2, we are really to wait for higher-energy experiments to verify (2.27).

## III. DISCUSSION

We note that in the present model for diffractive dissociation there is effectively incoherent scattering for one hadron and diffraction scattering for the other hadron. The hadron undergoing incoherent scattering with impulse approximation finally gets converted to the observed hadrons. This aspect here in fact becomes very similar to deep-inelastic lepton-hadron scattering.

It is interesting to observe that the present model generates the same expression as triple-Pomeron interaction in the high-energy limit with additional information regarding the slope parameter, as earlier observed by Ravndal.[10] The model appears to yield a new understanding of diffractive dissociation processes in the context of quark-parton models. From Fig. 2 it is obvious that the present model as expected is progressively valid at high energies, and in this context the disagreement in the other figures may be regarded as temporary low-energy effects. We may also conclude the same from the analysis of these events in the context of triple-Regge phenomenology,[21] where other trajectories besides the Pomeron need be included at the present energies, and such energy-dependent effects will not be simulated by the interaction (2.1).

For incoherent scattering of a proton we may imagine that the proton breaks up into a system of quarks and gluons,[11,12] and this system gets converted to the observed hadrons with unit probability. However, we note that with low enough $M^2$ given by (2.11), this system *cannot* go over to other hadrons just from kinematic considerations. This is an obvious known limitation of quark-parton models which requires in particular that $M^2$ be adequately large. We attempt to take into account this limitation in a smooth manner by associating a probability $p(M)$ with the disturbed hadron[11] of mass $M$. This $p(M)$ expresses the probability that the above system of mass $M$ gets converted to hadrons in a self-consistent manner. We expect that $p(M)$ is small when $M$ is near threshold and that $p(M)$ is 1 when $M$ is adequately large. To consider $p(M)$ for arbitrary $M$, let us in particular consider a disturbed proton $p^*$ obtained after incoherent scattering of a quark of the proton, which with asymptotic freedom[26] we assume consists of free quarks and gluons. This assumption, however, is to be further consistent with color confinement. We take this as equivalent to the assumption that free quarks and gluons can not exist separated by a distance $r_c$ or larger. To correlate these mutually contradictory assumptions of "freedom" with "confinement" we introduce a mean time for hadronization of $p^*$ as the system of quarks and gluons where this inverse lifetime or width is taken as

$$\Gamma(M)=\Gamma(p^*\to \text{hadrons})\ . \tag{3.1}$$

We assume for simplicity that $p^*$ breaks up into two hadrons in the neighborhood of threshold. Then the probability for hadronization of $p^*$ within a distance of $r_c$ is given by

$$p(M)=\{1-\exp[-\Gamma(M)\,r_c/2]\}\,, \tag{3.2}$$

where naturally we have chosen the rest frame of $p^*$ as the relevant frame of reference and have assumed that $p^*$ must hadronize within time $r_c/2$. As stated earlier, we take $p(M)$ given by (3.2) as the probability that the assumption of incoherent scattering for a given $M$ be valid. Hence, e.g., we shall now replace Eq. (2.15) by

$$\frac{d\sigma}{dt\,dM^2}\,(pp\to pX)=\frac{27\,f_V^{\,2}}{\pi}\left[1-\exp\left(-\Gamma(M)\,\frac{r_c}{2}\right)\right]\times\frac{F_{2p}\,x_F}{M^2-m^2}\,\exp(\tfrac{1}{3}R_p^{\,2}t)\ . \tag{3.3}$$

Now, Eq. (3.3) can be extrapolated down to small $M^2$ for the diffractive dissociation processes. When $M$ is small, $\Gamma(M)$ in (3.1) is small and thus $p(M)$ in (3.2) will be small, which will suppress the cross section in (3.3). When $M$ is large we may expect that (2.15) and (3.3) will be the same, since $\Gamma(M)$ will be large.

To give a heuristic estimate of $\Gamma(M)$ let us assume that $p^*$ is merely a heavy proton of mass $M$ with, e.g., a coupling described by

$$\sqrt{2}\,G\bar{n}(x)\gamma_5 p^*(x)[\pi^+(x)]^\dagger \tag{3.4}$$

for the decay of $p^*$ to a $\pi^+$ and a neutron. Then using isospin invariance we obtain that

$$\Gamma(p^*\to N\pi)=\frac{G^2}{4\pi}\,\frac{3p^3}{M(p^0+m)}\ , \tag{3.5}$$

where $p$ is the c.m. momentum for the pion-nucleon system. We take $G^2/4\pi=14.6$ as for the pion-nucleon system, and further,[27] $r_c=1$ fm for the domain of color confinement. With (3.3) we now draw curve II in Fig.1, which is expected to be valid for small $M^2$ also. We note that for $M^2\lesssim 4$ $\text{GeV}^2$, the data points lie above the plotted curve II. This excess cross section is interpreted as being due to resonance production, as has been stated earlier. This interpretation has the further support that the slope parameter in this region is too high, as is expected, since the resonance production process is similar to that of diffraction scattering.[4,10]

The region of small $M^2$ has been carefully investigated recently.[28] We have plotted in Fig. 6 these data, where, in the present model, we *separate* the cross section for diffractive dissociation and that for resonance production. The lower curve A is the same as curve II of Fig. 1, corresponding to diffractive dissociation. The balance is attributed to resonance production.

We note that in this region the only resonance present is $N(1400)$, which is with the same quan-

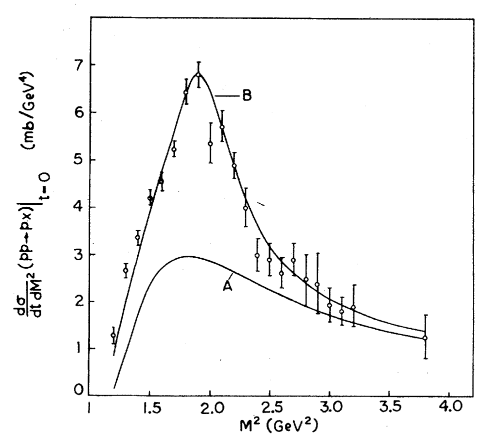


FIG. 6. $(d\sigma/dt\,dM^2)\ (pp\to pX)|_{t=0}$ vs $M^2$ for small $M^2$. Curve A is due to diffractive dissociation only as per (3.3). Curve B is the combined effect of resonance $N(1400)$ as per (3.6) along with the contribution from diffractive dissociation. The experimental points are from Ref. 28 (Table I, high-energy limit).

tum numbers[29] as the proton, and hence we may presuppose that this is the only resonance which is produced. For the production of this resonance, we assume a Breit-Wigner form with

$$\frac{d\sigma}{dt\,dM^2}=\frac{A'\exp(b_0't)}{2M\left[(M-m_R)^2+(\Gamma/2)^2\right]}, \qquad (3.6)$$

where we have taken $m_R=1.39$ GeV, $\Gamma=0.24$ GeV, and have adjusted the constant $A'$. In Fig. 6 curve B we have then plotted the combined effect of (3.3) and (3.6) and compared the results with experiments. The results are in good agreement with the experiments, which demonstrates that $N(1400)$ may be the only resonance produced here. Taking $b_0'\simeq 15$ GeV$^{-2}$, we also obtain from (3.6) that $\sigma_t(pp\to N(1400)p)\simeq 0.53$ mb. This result is the same as in some experiments.[30] Since, as per the results of the present model, diffractive dissociation will always be associated with a slope of 6 GeV$^{-2}$, it may be worthwhile to attribute two slopes in this region, one for diffractive dissociation and the other for resonance production, for an analysis of experimental points, with different proportions of both. In this case the large errors in Ref. 28 for the slope may no longer be there.

While considering limitations for the incoherent scattering, it is clear that Eqs. (3.1) and (3.2) are parametrizations of a very complicated process, which is nonperturbative, and where no solutions are known to exist.[31] Besides the heuristic reasons stated earlier, we may also tend to believe in the nature of the above explanations in view of the internally consistent picture regarding resonance production and diffractive dissociation, particularly in Fig. 6 as well as in Fig. 1. A similar analysis regarding diffractive dissociation of the meson for very small $M^2$ has not been carried out here because of the absence of any clear experimental picture.

We can thus obtain at a phenomenological level an understanding of purely diffractive[4] and diffractive dissociation processes at high energies with a current-current interaction in quark space in a field-theoretic quark model,[2] correlating the above processes to the form factors and to deep-inelastic lepton-hadron processes.[11] We may note that our results constitute the high-energy limit, whereas the experimental results are generally energy dependent, as may be quite clear in Fig. 2. Further, the disagreements of the results for high $M^2$, say $M^2 \gtrsim 30$ GeV$^2$, may be noted. This indicates that the phenomenological Hamiltonian (2.1) with factorizability as in (2.3) is no longer adequate. This may be due to more complicated dynamics including gluon production[12] in the primary collision in quark space and subsequent fragmentation.[13]

We note that the results we have derived with the quark model are similar to that of triple-Pomeron interaction, and also sometimes go beyond the Regge approximation with Pomeron exchange in the context of quark-parton ideas. We may add that the current corresponds to a spin-1 exchange with vacuum quantum numbers and thus, the correspondence with the Pomeron is not really unexpected. The present ideas, however, also throw new light on the limitations of quark-parton models as in (3.2), in addition to indicating possible dynamics for the diffractive dissociation processes in the context of the quark model.

## ACKNOWLEDGMENTS

The authors are thankful to T. Pradhan, J. C. Pati, A. Khare, L. Maharana, K. Biswal, and S. K. Mishra for discussions. One of the authors (A.R.P.) is indebted to the University Grants Commission for financial assistance.

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