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Photoproduction of vector mesons at high energies and Compton scattering by nucleons

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We consider photoproduction of vector mesons and Compton scattering at high energies. Coupling of the electromagnetic current to neutral vector mesons as given theoretically in quark models and experimentally in storage-ring experiments is assumed. This gives a reasonable theory for ρ, ω, ϕ, ψ, and ψ' photoproduction processes in a unified manner with appropriate energy dependence for $(d\sigma/dt)_0$ for ϕ, ψ, and ψ' as well as for the slope parameter while approaching the high-energy limit. The correct total Compton-scattering cross section by nucleons in the high-energy limit is obtained. It appears that one is able to understand in a qualitative manner the success of vector-dominance models in the context of the quark model, without assuming vector dominance in a conventional form and without any field-current identity.

I. INTRODUCTION

It has been known for quite some time that a photon very often behaves like a hadron while interacting with hadrons. This behavior is beautifully explained in field theory through the vector-dominance model[1] originally proposed by Sakurai. This hypothesis includes *all* interactions of photons with hadrons. However, we believe that hadrons are quark and antiquark composites, and since quarks are charged, photons should interact with the quarks. In such a picture, one should have at least a qualitative understanding of vector dominance or a similar model as a derived hypothesis rather than a primary assumption for the interaction of photons with hadrons. The success of vector dominance indicates that there should be a coupling of vector mesons with photons. Such a picture was imagined by Gottfried and Yennie[2] while considering shadowing phenomena for photon absorption by nuclei. Here a photon–vector-meson coupling was assumed, the nucleus was replaced by a complex potential, and with a potential scattering in the eikonal approximation for high energies, the behavior of the photon while interacting with nucleons of nuclei could be reasonably understood. Effectively, except for the photon–vector-meson coupling, vector dominance was not utilized.

Photon–vector-meson coupling can be easily understood in the context of a quark model since $\langle \mathrm{vac}| J^{\mu}(0)| V\rangle \neq 0$, and is related to the wave function of the vector meson at the origin, as was first noted by Van Royan and Weisskopf.[3] Some corrections to this relation were noted by one of the authors in a specific four-component quark model which takes into account quark motion inside hadrons.[4] A specific theory of Lorentz boosting was also given to describe hadrons in motion[5] in a way similar to the representation of a Lorentz group through the rotation group as the little group.[6] In this version of the relativistic quark model,[4,5] it appears that we may have a reasonable quantitative generation of strong interactions consistent with the Okubo-Zweig-Iizuka rule[7] from only static-quark-model parameters. The results for radiative decays of vector mesons[8] obtained previously by us also seem to further justify the specific theory of Lorentz boosting proposed.[5] It was also conjectured in Ref. 5 that merely the nonvanishing of $\langle \mathrm{vac}| J^{\mu}(0)| V\rangle$ may be adequate to understand the results of vector-meson dominance[1] (VMD) in the field-theoretic language of photoproduction and Compton scattering by nucleons. We examine this hypothesis in this paper and find that in fact the photon is likely to behave like a hadron with an exchange of vector mesons at high energies in the relativistic quark model as proposed by one of the authors.[5]

II. GENERAL THEORY

It is generally believed that quarks are permanently confined in hadrons.[9] We shall make this assumption here. In this context, it was conjectured that the quark-field operators span an unphysical vector space.[5] The physical vector space, which we should take in our perturbation theories, is that of hadrons, obtained as eigenstates of the Hamiltonian as in Ref. 4. We do not have quark propagators; but we can have hadronic intermediate states in our perturbation theory, which will simulate hadronic propagators. We shall first develop such a perturbation theory in the lowest orders. Although this result is trivial, we shall state it here for the sake of completeness. Let us have

$$V(t) = \int \mathcal{H}_I(x) d^3x\,. \tag{2.1}$$

In the above equation, $\mathcal{H}_I(x)$ is an invariant local field operator, including quark-field operators which will generate hadronic interactions, and *does not* contain hadronic-field operators. We can possibly define effective hadronic-field operators, but we do not do so since this will take our attention away from the hadrons being composite, and we imagine all hadronic interactions to be *generated* through quark interactions. Clearly, we are to take $\psi_Q^\xi(x)$ defined in Ref. 5 in $_I(x)$ to make the theory independent of any Lorentz frame of hadrons.[5] The S matrix is given as[10]

$$S = I + \sum_{n=1}^{\infty} S_n ,$$

where

$$S_n = (-i)^n \int_{-\infty}^{\infty} dt_1 \int_{-\infty}^{t_1} dt_2 \cdots \int_{-\infty}^{t_{n-1}} dt_n V(t_1) \cdots V(t_n) . \tag{2.2}$$

Taking *space* and time translational invariance, and introducing *hadronic*, leptonic, and photon intermediate states in (2.2), we then obtain

$$\langle f|S|i\rangle = \delta_4(P_f - P_i) M_{fi} , \tag{2.3}$$

where the nth-order contribution $M_{fi}^{(n)}$ is given as

$$M_{fi}^{(n)} = (-i)(2\pi)^{3n+1} \sum_{i_1,\ldots,i_{n-1}} \langle f|\mathcal{H}_I(0)|i_1\rangle\langle i_1|\mathcal{H}_I(0)|i_2\rangle \cdots \langle i_{n-1}|\mathcal{H}_I(0)|i\rangle \prod_{k=1}^{n-1} (E_i - E_k + i\epsilon)^{-1} . \tag{2.4}$$

In (2.4), E_k is the energy of the state i_k. The appropriate singularity in the energy denominators has been taken. Also, for $k = 1, \ldots, n-1$, $\vec{P}_k = \vec{P}_i = \vec{P}_f$, i.e., total momentum is conserved in the intermediate states, and summation over i_k, $k = 1, \ldots, n-1$, indicates summation for the remaining degrees of freedom. We shall employ (2.4) for our calculations in the lowest order for the purpose of photoproduction of vector mesons and Compton scattering. In our field-theoretic version of composite hadrons, the old-fashioned perturbation-theory result (2.4) converted to field-theoretic assumptions becomes very useful.

III. PHOTOPRODUCTION OF VECTOR MESONS BY NUCLEONS

We shall now consider photoproduction of vector mesons, using (2.4). As usual, we shall assume for this purpose that the amplitude of $VN \to VN$, where V is a vector meson and N is the nucleon, is known. Specifically, we shall take that in (2.1),

$$\mathcal{H}_I(x) = \mathcal{H}_{\rm em}(x) + \mathcal{V}(x) . \tag{3.1}$$

In the above,

$$\mathcal{H}_{\rm em}(x) = eJ^\mu(x) A_\mu(x) , \tag{3.2}$$

where $J^\mu(x)$ is the electromagnetic current[5] expressed in terms of quark-field operators $\psi_Q^\xi(x)$. Further, $\mathcal{V}(x)$ is the invariant effective quark-field-operator Hamiltonian which correctly describes $VN \to V'N'$. Thus matrix elements of $\mathcal{V}(0)$ will be assumed to be determined from $(d\sigma/dt)(VN \to V'N')$. The universality of diffraction scattering at high energies gives us reasonable estimates for this expression. We note that regarding this expression we are not doing any better than conventional theories.

We first note that $\mathcal{H}_{\rm em}(x)$ *alone* in the second order will give rise to vector-meson productions through the *photon absorption* process

$$\gamma N \to B \to V'N' , \tag{3.3}$$

where B is an intermediate baryon state. However, in the quark model, for the process (3.3) there will be *two spectator* quarks of the nucleon and B states, with one quark absorbing the photon. With the quark-field operators describing this process, the momentum of the spectator quarks will remain unaltered, whereas the momentum of the quark which absorbs the photon will become high.[5] Thus the approximate form of the harmonic-oscillator wave function of N and B will completely suppress this matrix element.[11] Again, when we use the c.m. frame of reference, the energy denominator $p^0 + |\vec{k}| - p_B^0$ in (2.4) will further suppress this contribution. Hence we conclude that for hadrons, direct absorption of high-energy photons by the hadron will be highly suppressed. This is both due to the composite structure of the hadrons described approximately by harmonic-oscillator wave functions as well as due to large denominators in the perturbation theory at high energies. We may remark that the above suppression will also work for any absorption vertex, e.g., meson absorption through the pair-annihilation component of the Dirac Hamiltonian.[4,5]

We should also recognize that the above comments are of a qualitative nature. Particularly, all the available baryon states B in (3.3) will contribute, giving rise to an infinite summation, possibly with small denominators in this contribution for high-mass baryons. In view of our ignorance of such contributions, it is impossible to assess their importance. However, we shall "guess" that for the reasons mentioned in the last paragraph, the contribution from (3.3) will be negligible. As our subsequent calculations show, this enables us to understand the origin of vector dominance in a qualitative manner.

We shall thus modify the conventional vector-dominance hypothesis as a combination of the two interaction terms in (3.1) with an exchange of vector meson. For this purpose we shall calculate the effect of $\mathcal{V}(x_1)\mathcal{H}_{em}(x_2)$ in the second order in (2.4), and note that when we take the universal cross sections for diffraction scattering, there is no suppression of this contribution. In fact, for reasons stated earlier, it appears that this contribution will dominate over direct absorption as envisaged in (3.3).

We take $(\vec{k},\vec{e})$ and $(-\vec{k},\gamma)$ to describe momenta and polarizations of the initial photon and nucleon and $(\vec{k}',\lambda')$ and $(-\vec{k}',r')$ to describe the same for the outgoing vector meson and nucleon. Thus we get in the second order, by (2.4),

$$M_{fi}=(-ie)(2\pi)^7(k^0-k_V^0)^{-1}\times\sum_\lambda\langle-\vec{k}'r';\vec{k}'\lambda'|\mathcal{V}(0)|-\vec{k}r;\vec{k}_V\lambda\rangle\times\langle\vec{k}_V\lambda|J^\mu(0)A_\mu(0)|\vec{k}\vec{e}\rangle. \tag{3.4}$$

In the above equation

$$\vec{k}_V=\vec{k}\text{ and }k_V^0=(m_V^2+\vec{k}^2)^{1/2}. \tag{3.5}$$

We now note that[5]

$$\begin{aligned}\langle\vec{k}_V\lambda|J^\mu(0)|\text{vac}\rangle&=\left(\frac{m_V}{k_V^0}\right)^{1/2}\langle\vec{0}\lambda|U^{-1}[L(k_V)]J^\mu(0)|\text{vac}\rangle\\&=\left(\frac{m_V}{k_V^0}\right)^{1/2}L_{\mu\nu}(k_V)\langle\vec{0}\lambda|J^\nu(0)|\text{vac}\rangle\\&=\frac{m_V^2}{f_V}(2\pi)^{-3/2}\frac{1}{(2k_V^0)^{1/2}}\epsilon^\mu(k_V,\lambda).\end{aligned} \tag{3.6}$$

In (3.6) we have used[3,4]

$$\langle\vec{0}\lambda|J^\nu(0)|\text{vac}\rangle=\frac{m_V^2}{f_V}(2\pi)^{-3/2}\frac{1}{(2m_V)^{1/2}}\delta_{\nu\lambda}, \tag{3.7}$$

and[12]

$$L_{\mu\lambda}(k_V)=\epsilon^\mu(k_V,\lambda). \tag{3.8}$$

Clearly, $L(k_V)$ is the Lorentz-boosting matrix taken earlier.[5,12] Although notations in (3.6) have the conventional form for field-current identity,[13] no such identity has been assumed and in fact we have no vector-meson field operator. Although we can possibly define such an object, we do not do so since it may not correctly describe the interaction of vector mesons except in an *ad hoc* manner. f_V is experimentally known from the coupling of vector mesons to the e^+e^- channel. We thus obtain from (3.4) and (3.6)

$$M_{fi}(\gamma N\to V'N')=(-ie)(2\pi)^4(k_V^0-k^0)^{-1}\times\sum_\lambda\vec{\epsilon}(\vec{k}_V,\lambda)\cdot\vec{e}\,\frac{1}{(2k_V^0 2k^0)^{1/2}}\frac{m_V^2}{f_V}\langle\vec{k}',\lambda';-\vec{k}',r'|\mathcal{V}(0)|\vec{k}_V,\lambda;-\vec{k},r\rangle\Big|_{\vec{k}_V=\vec{k}}. \tag{3.9}$$

We shall determine the matrix element on the right-hand side of (3.9) by speculating on $VN\to V'N'$ diffraction scattering. At the outset we note that the matrix element

$$\langle\vec{k}',\lambda';-\vec{k}',r'|\mathcal{V}(0)|\vec{k}_V,\lambda;-\vec{k},r\rangle$$

conserves momentum, but *does not* conserve energy. We define s and t as the usual variables associated with the process $\gamma N\to V'N'$, and further take

$$s_0=(k_V^0+p^0)^2,\quad t_0=(k_V-k')^2, \tag{3.10}$$

where $p^0=(m^2+\vec{k}^2)^{1/2}$, and m is the mass of the nucleon. We now write from general invariance considerations and from the fact that at high energies there is hardly any spin or unitary-spin flip

$$\langle\vec{k}',\lambda';-\vec{k}',r'|\mathcal{V}(0)|\vec{k}_V,\lambda;-\vec{k},r\rangle=\delta_{\lambda\lambda'}\delta_{r'r}\left(\frac{m_V}{k_V'^0}\right)^{1/2}\left(\frac{m_V}{k_V^0}\right)^{1/2}\left(\frac{m}{p'^0}\right)^{1/2}\left(\frac{m}{p^0}\right)^{1/2}\times A(s,t;s_0,t_0), \tag{3.11}$$

where $p'^0=(m^2+\vec{k}'^2)^{1/2}$. Equation (3.9) then simplifies to

$$M_{fi}(\gamma N\to V'N')=(-ie)(2\pi)^4(k_V^0-k^0)^{-1}\vec{e}_{\lambda'}\frac{\delta_{r,r'}}{(2k_V^0 2k^0)^{1/2}}\times\frac{m_V^2}{f_V}\left(\frac{m_V^2m^2}{k_V'^0k_V^0p'^0p^0}\right)^{1/2}A(s,t;s_0,t_0). \tag{3.12}$$

We now note that with $\mathcal{V}(x)$ in the first order and with (3.11),

$$M_{fi}(VN\to V'N')=(-i)(2\pi)^4\delta_{r'r}\delta_{\lambda'\lambda}\left(\frac{m_Vm}{k_V'^0p'^0}\right)A(s,t'), \tag{3.13}$$

where $-t'$ is the momentum transfer squared for the above s for $VN \to V'N'$. Explicitly in (3.13), we take the momenta as $(\vec{k}_0, -\vec{k}_0) \to (\vec{k}', -\vec{k}')$, where $\vec{k}_0 = (\kappa'/\kappa)\vec{k}$, so that (3.13) is physically admissible with known $V'N'$. We note that then

$$t' = -2\kappa'^2(1-\cos\theta)\,, \tag{3.14}$$

where

$$\vec{k}\cdot\vec{k}' = \kappa\kappa'\cos\theta\,. \tag{3.15}$$

As compared to the above,

$$t = (k_V' - k)^2 = (k_V'^0 - k^0)^2 - 2\kappa\kappa'(1-\cos\theta) - (\kappa-\kappa')^2\,. \tag{3.16}$$

We now note that[14] from (3.13)

$$\frac{d\sigma}{dt'}(VN \to V'N') = \frac{4\pi^3}{\kappa'^2(p'^0 + k_V'^0)^2} m^2 m_V{}^2 (2\pi)^8 \times |A(s,t')|^2\,. \tag{3.17}$$

On the other hand, from (3.12) we get[14]

$$\frac{d\sigma}{dt}(\gamma N \to V'N') = \frac{4\pi^3}{\kappa^2(p^0+k^0)^2}\frac{4\pi\alpha}{f_V{}^2}\frac{m_V{}^4}{4k_V^{02}(k_V^0-k^0)^2} \times m^2 m_V{}^2 (2\pi)^8 |A(s,t;s_0,t_0)|^2\,. \tag{3.18}$$

We now *assume* that

$$A(s,t') \simeq A(s,t;s_0,t_0)\,. \tag{3.19}$$

We can then correlate (3.17) and (3.18), and thus obtain

$$\frac{d\sigma}{dt}(\gamma N \to V'N') = \frac{4\pi\alpha}{f_V{}^2}\frac{\kappa'^2}{\kappa^2}\frac{m_V{}^4}{4k_V^{02}(k_V^0-k^0)^2} \times \frac{d\sigma}{dt'}(VN \to V'N')\,. \tag{3.20}$$

We first note that for very-high-energy photons $\kappa' \simeq \kappa$, and thus we obtain the standard vector-dominance-model result that for $\kappa \to \infty$,

$$\frac{d\sigma}{dt}(\gamma N \to V'N') = \frac{4\pi\alpha}{f_V{}^2}\frac{d\sigma}{dt}(VN \to V'N')\,. \tag{3.21}$$

However, at intermediate energies, there are basic energy-dependent factors. First, let us parametrize

$$\frac{d\sigma}{dt'}(VN \to V'N') = \left(\frac{d\sigma}{dt}\right)_0 \exp(bt')\,. \tag{3.22}$$

We now note from (3.14) and (3.16) that

$$t' = \frac{\kappa'}{\kappa}(t - t_{\min})\,, \tag{3.23}$$

where

$$t_{\min} = (k_V'^0 - k^0)^2 - (\kappa' - \kappa)^2\,. \tag{3.24}$$

We then obtain from (3.20)

$$\frac{d\sigma}{dt}(\gamma N \to V'N') = \frac{4\pi\alpha}{f_V{}^2}\frac{\kappa'^2}{\kappa^2}\frac{m_V{}^4}{(k_V^0-k^0)^2 4k_V^{02}} \times \left(\frac{d\sigma}{dt}\right)_0 (VN \to V'N') \times \exp\left(\frac{\kappa'}{\kappa}b(t-t_{\min})\right). \tag{3.25}$$

Now, comparing (3.22) and (3.25), we note that even when b is constant, the slope parameter in (3.25) will be smaller than b and will be energy dependent. Further, with laboratory photon energy E_γ, there will be an energy-dependent suppression factor for $(d\sigma/dt)_0(\gamma N \to V'N')$ given by

$$f_V(E_\gamma) = \frac{\kappa'^2}{\kappa^2}\frac{m_V{}^4}{4k_V^{02}(k_V^0-k^0)^2}\,. \tag{3.26}$$

Clearly, we have

$$|\vec{k}| = \kappa = k^0 = \frac{(m)^{1/2}}{(m+2E_\gamma)^{1/2}}E_\gamma \tag{3.27}$$

and

$$|\vec{k}'| = \kappa' = \left(\frac{\lambda'}{4s}\right)^{1/2}, \tag{3.28}$$

where[15]

$$\lambda' = (s - m^2 - m_V{}^2)^2 - 4m^2 m_V{}^2\,. \tag{3.29}$$

A. ρ^0 and ω production

Let us compare our results with experiments.[16] We note that

$$\frac{d\sigma}{dt}(\gamma N \to V'N') = A\exp[B(t-t_{\min})]\,, \tag{3.30}$$

where

$$A(E_\gamma) = \left(\frac{d\sigma}{dt}\right)_0 (\gamma N \to V'N') = \frac{4\pi\alpha}{f_V{}^2}f(E_\gamma)\left(\frac{d\sigma}{dt}\right)_0 (VN \to V'N') \tag{3.31}$$

and

$$B = \frac{\kappa'}{\kappa}b\,, \tag{3.32}$$

with the defining equation (3.22) for b. We shall assume as usual amplitudes are pure imaginary, the photoproduction process is mainly diffractive, and that $\sigma_{\text{tot}}(VN \to VN)$ and b are independent of energy. We then obtain

$$\sigma_t(\gamma N \to V'N') = A/B \tag{3.33}$$

and

$$\sigma_t(VN \to VN) = \left[\frac{4f_V{}^2 A}{\alpha f(E_\gamma)}\right]^{1/2} = \left[\frac{16\pi\alpha m_V}{3\Gamma_V}\frac{A}{f(E_\gamma)}\right]^{1/2}, \tag{3.34}$$

TABLE I. ρ^0 photoproduction. The input constitutes 9.3-GeV photoproduction data (Ref. 17).

E_γ (GeV)	A (μb/GeV^2)	B (GeV^{-2})	σ_t ($\gamma N \to \rho^0 N$) (μb)
2.8	61.1	6.4	9.5
4.7	78.3	7.0	11.2
9.3	95	7.3	13
∞	109.5	7.6	14.4

TABLE II. ϕ photoproduction. Inputs are σ_t (ϕN) = 10.8 mb and b (ϕN) = 6.5 GeV^{-2}.

E_γ (GeV)	A (μb/GeV^2)	B (GeV^{-2})	σ_t ($\gamma N \to \phi N$) (μb)
2	0.662	3.34	0.20
3	1.02	3.94	0.26
4	1.99	5.35	0.37
5	2.29	5.63	0.41
6	2.50	5.80	0.43
7	2.65	5.91	0.45
8	2.77	6.0	0.46
9	2.86	6.05	0.47
10	2.94	6.1	0.48
∞	3.69	6.5	0.57

where $\Gamma_V = \Gamma(V \to e^+ e^-)$, and the width is calculated without any corrections for the finite width of the vector mesons.[14] Obviously (3.34) is assumed to be constant in our calculations, which gives the energy dependence of A. Using these assumptions, and with the 9.3-GeV data as input,[17] we give the calculated results in Table I. We note that the values of B are consistent with the experimental results, and this explains the change in slope with the asymptotic value of 7.6 GeV^{-2} as the slope for ρN scattering, in general agreement with meson-baryon diffraction scattering at high energies. However, the values of A in Ref. 16 have just the *opposite* energy dependence than what we have taken. This may be due to the fact that scattering may not be via the channel we consider—e.g. the one-pion-exchange term[17,18] is known to have about a 50% contribution at 2.8 GeV for ω photoproduction. Also, assumption (3.19) may not be valid at low energies, where kinematic corrections can be relevant. However, in this context we may note that Bapu *et al.*[19] obtain

$$\left(\frac{d\sigma}{dt}\right)_0 (\gamma p \to \rho^0 p) \simeq 50 \pm 15 \ \mu\text{b/GeV}^2$$

at 3 GeV for deuteron targets, with dissociation of the deuteron, which is in agreement with predictions of Table I, along with the slope parameter, which, however, has too much or an error.

We may note from Table I, that upon using (3.34) and including finite-width corrections,[14] and with[20] $\Gamma_\rho = 6.45$ keV, we obtain

$$\sigma_t(\rho N) = 26.5 \text{ mb}. \tag{3.35}$$

For photoproduction of ω we expect a similar behavior as above if we only take the diffractive part, and the results are consistent with this after the one-pion-exchange contribution is separated at low energies.[17,18] When we take[20] $\Gamma_\omega = 0.76$ keV, and[17] $A = 13.7$ μb/GeV^2 at 9.3 GeV, we again obtain, from (3.34),

$$\sigma_t(\omega N) = 26 \text{ mb}, \tag{3.36}$$

which agrees with (3.35). We get the ratio f_ρ^2/f_ω^2, the same as is expected from broken SU(6) symmetry,[21] and the total cross sections in (3.35) and (3.36) are as per the predictions of the quark model for meson baryon scattering.

B. ϕ production

For ϕ-meson photoproduction, the energy dependence of A and B in (3.31) and (3.32) is more pronounced theoretically (see Table II) and the experimental data are also available.[22,23] We plot the data of Ref. 22 in Fig. 1, along with the predicted curve. As in (3.35), with[20] $\Gamma_\phi = 1.31$ keV and finite-width corrections,[14] we find that this curve corresponds to

$$\sigma_t(\phi N) = 10.8 \text{ mb}. \tag{3.37}$$

The agreement with the quark-model prediction[24] of $\sigma_t(\phi N) = 13$ mb is reasonable. We thus note that with the energy dependence of $f(E_\gamma)$, $f_\phi^2/4\pi$ as derived from storage-ring experiments and from the photoproduction data becomes mutually

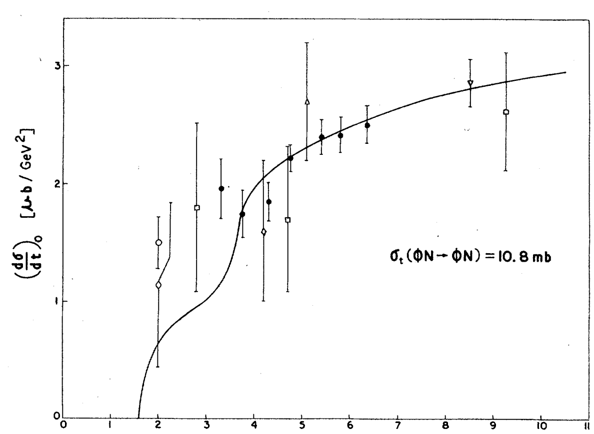

FIG. 1. $(d\sigma/dt)_0$ for ϕ production versus E_γ along with the predicted curve. Data taken from Ref. 22, Fig. 8.

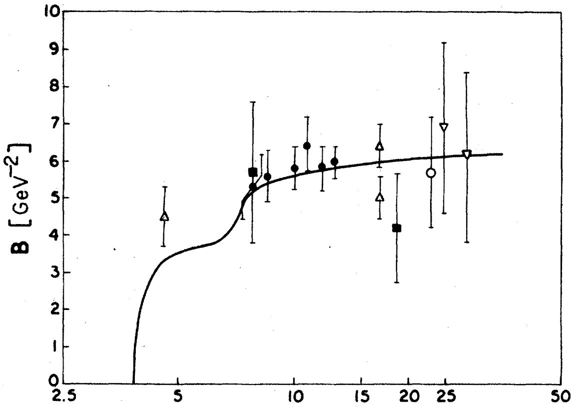

FIG. 2. Slope parameter B for $|t| \leq 0.4$ GeV^2, along with the predicted curve. Data taken from Ref. 23, Fig. 2. 3.

consistent, and the anomaly regarding this disappears.[25]

We also plot the data for B, the slope parameter, against s in Fig. 2, with the predicted curve where the input is $b(\varphi N \to \varphi N) = 6.5$ GeV^{-2}. We notice from Figs. 1 and 2 that the agreement is quite good for $E_\gamma \gtrsim 4$ GeV, i.e., $s \gtrsim 8$ GeV^2. We further note that even at $E_\gamma = 10$ GeV, the asymptotic form is not reached.

Since at these energies there is considerable energy dependence, and with the present error bars there is quantitative agreement, it may be worthwhile to conduct experiments with a smaller energy spread in E_γ.

C. ψ and ψ' production

As earlier, we calculate A and B as a function of E_γ and compare these with experimental results[26] in Table III. $\sigma_t(\psi N) = 1$ mb has been taken as the input, along with the assumption that $b(\psi N \to \psi N) = 4$ GeV^{-2} as obtained for experiments at[27] $E_\gamma = 100$ GeV.

We notice that obviously the agreement is reasonable, except for very-low-energy data. In Table III we have used[20] $\Gamma_\psi = 4.69$ keV. If we further take[20] $\Gamma_{\psi'} = 2.05$ keV, we then obtain for $E_\gamma = 21$ GeV

$$\frac{(d\sigma/dt)_0(\psi)}{(d\sigma/dt)_0(\psi')} = \frac{\Gamma_\psi}{\Gamma_{\psi'}} \frac{m_{\psi'}}{m_\psi} \frac{f_\psi(21\text{ GeV})}{f_{\psi'}(21\text{ GeV})} \simeq 4.1 . \tag{3.38}$$

In (3.38), we have assumed that $\sigma_t(\psi N) = \sigma_t(\psi' N)$. The experimental value[28] for (3.38) is 6.8 ± 2.4, which is quite satisfactory considering the margin of error involved in the experiments and the parameters. We also note that for ψ the slope parameter as available[26] at 19 GeV is $B = 2.9 \pm 0.3$ GeV^{-2} compared to the predicted value 3.0 GeV^{-2} in Table III, which shows that diffraction slopes for φN and ψN scattering are not the same. We thus again explain the change of slope at[26] 19 GeV and at[27] 100 GeV.

We have plotted A against s with the experimental points[26,28] against the theoretical curve in Fig. 3. The agreement is reasonable.

IV. COMPTON SCATTERING BY NUCLEONS

We shall take the same Hamiltonian as in (3.1). The first term on the right-hand side of (3.1) will yield Compton scattering with the absorption of a high-energy photon by a hadron, which will correspond to a conventional diagram for Compton scattering. However, as discussed in the last section, this process will be highly suppressed due to the composite nature of the hadrons as well as the large energy denominators. On the other hand, we shall see that no such suppression will operate for the third-order calculations in (2.4) with an exchange of vector mesons. As we shall see this yields the results of the vector-dominance model of field theory, with kinematic corrections as in the last section.

TABLE III. ψ photoproduction. Inputs are $\sigma_t(\psi N) = 1$ mb and $B = (\kappa'/\kappa)\, b(\psi N) = 4$ GeV^{-2} from Ref. 27 at $E_\gamma = 100$ GeV, and experimental points are from Refs. 26 and 28.

E_γ (GeV)	A (nb/GeV^2)	B (GeV^{-2})	$\sigma_t(\gamma N \to \psi N)$ (nb)	A (expt.) (nb/GeV^2)
9	1.3	1.1	1.2	...
13	6.7	2.3	2.91	3.8 ± 0.8
15	8.8	2.6	3.38	6.8 ± 2
16	9.7	2.7	3.59	8.2 ± 1.1
17	10.5	2.8	3.75	10.8 ± 1
19	12.0	3.0	4.0	13.5 ± 3
21	13.3	3.1	4.3	14.6 ± 1.2
100	27.2	4	6.8	40 ± 13
∞	33.2	4.2	7.90	...

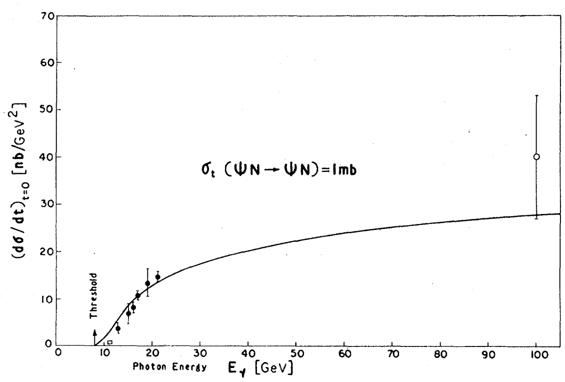

FIG. 3. $(d\sigma/dt)_0$ for ψ production versus E_γ along with the predicted curve. Data taken from Refs. 26 and 28.

The third-order contribution in (2.4) with (3.1) will come from the product of the operator in the form

$$\mathcal{H}_{\rm em}(x_1)\mathcal{V}(x_2)\mathcal{H}_{\rm em}(x_3)\,. \tag{4.1}$$

The contribution will thus come from two admissible types of intermediate states for Compton scattering of nucleons:

$$N\gamma \rightarrow NV \rightarrow N'V' \rightarrow N'\gamma' \tag{4.2}$$

and

$$N\gamma \rightarrow N\gamma V'\gamma' \rightarrow N'\gamma V\gamma' \rightarrow N'\gamma'\,. \tag{4.3}$$

Let us take $(\vec{k}, \vec{e})$ and $(-\vec{k}, r)$ and $(\vec{k}', \vec{e}')$ and $(-\vec{k}', r')$ as the momenta and polarizations of the initial photon and nucleon and final photon and nucleon, respectively. We shall not consider the contribution from (4.3), since at high energies this will vanish due to the large energy denominators and, as can be seen, due to large momentum transfers in the corresponding strong process. This may be relevant for backward Compton scattering. Omitting this we now consider the third-order intermediate states as in (4.2). Then (2.4) yields the corresponding amplitude

$$M^a_{fi} = (-i)(2\pi)^{10} \sum_{V,\lambda,\lambda'} \frac{e^2}{(k_V^0 - k^0)^2} \langle \vec{k}', \vec{e}\,' | J^\mu(0) A_\mu(0) | \vec{k}_V', \lambda' \rangle \langle \vec{k}_V', \lambda'; -\vec{k}', r' | \mathcal{V}(0) | \vec{k}_V, \lambda; -\vec{k}, r \rangle$$
$$\times \langle \vec{k}_V, \lambda | J^\nu(0) A_\nu(0) | \vec{k}, \vec{e} \rangle\,. \tag{4.4}$$

In (4.4), clearly $\vec{k}_V = \vec{k}$ and $\vec{k}_V' = \vec{k}'$. λ and λ' are polarization states for the intermediate vector mesons. We now substitute (3.11) corresponding to diffraction vector-meson scattering. Since $|\vec{k}| = |\vec{k}'|$, clearly the question of energy nonconservation does not arise. Using s_0 as the energy in the vector-meson–nucleon channel, we write, corresponding to (3.11), with $p_0 = (\vec{k}^2 + m^2)^{1/2}$,

$$\langle \vec{k}_V', \lambda', -\vec{k}', r' | \mathcal{V}(0) | \vec{k}_V, \lambda; -\vec{k}, r \rangle$$
$$= \delta_{\lambda\lambda'}\, \delta_{rr'} \left(\frac{m_V}{k_V^0}\right)\left(\frac{m}{p^0}\right) A(s_0, t)\,. \tag{4.5}$$

In the above,

$$s_0 = [(m_V{}^2 + \vec{k}^2)^{1/2} + p^0]^2\,. \tag{4.6}$$

Proceeding now in a manner identical to that for the simplification of (3.4) and then (3.9), we thus obtain, from (4.4),

$$M^a_{fi} = (-i)(2\pi)^4 (\vec{e}\,' \cdot \vec{e})\delta_{r'r}$$
$$\times \sum_V \frac{e^2}{f_V{}^2} \frac{m_V{}^4}{(k_V^0 - k^0)^2 4k_V^{02}} \frac{m_V m}{k^0 p^0} A(s_0, t)\,. \tag{4.7}$$

However, now, corresponding to (3.17) we get

$$\frac{d\sigma}{dt}(VN \rightarrow V'N') = \frac{4\pi^3}{\kappa^2 (p^0 + k_V^0)^2} (m m_V)^2 (2\pi)^8$$
$$\times |A(s_0, t)|^2\,. \tag{4.8}$$

Also, (4.7) yields, ignoring all other contributions,

$$\frac{d\sigma}{dt}(\gamma N \rightarrow \gamma' N')$$
$$= \frac{4\pi^3}{\kappa^2 (p^0 + k^0)^2} (2\pi)^8 p^{02} |\vec{k}|^2$$
$$\times \left(\sum_V \frac{e^2}{f_V{}^2} \frac{m_V{}^4}{4k_V^{02}(k_V^0 - k^0)^2} \frac{m_V m}{k^0 p^0} A(s_0, t) \right)^2 . \tag{4.9}$$

Comparing (4.8) and (4.9), with the usual assumptions, we write, corresponding to (3.20),

$$\frac{d\sigma}{dt}(\gamma N \rightarrow \gamma' N') = \left[\sum_V \frac{4\pi\alpha}{f_V{}^2} \frac{m_V{}^4}{(k_V^0 - k^0)^2 4k_V^{02}} \frac{p^0 + k_V^0}{p^0 + k^0} \right.$$
$$\left. \times \left(\frac{d\sigma}{dt}(Vp \rightarrow Vp) \right)^{1/2} \right]^2 . \tag{4.10}$$

The high-energy limit of (4.10) becomes

$$\frac{d\sigma}{dt}(\gamma p \rightarrow \gamma' p') = \left[\sum_V \frac{4\pi\alpha}{f_V{}^2} \left(\frac{d\sigma}{dt}(VN \rightarrow V'N') \right)^{1/2} \right]^2 , \tag{4.11}$$

which is the conventional result of the vector-dominance model, corresponding to (3.21) for vector-meson photoproduction. We notice that the energy dependence in (3.20) and (4.10) is not the same.

From (4.10) we easily obtain that at high energies,

$$\sigma_t(\gamma N)=\sum_V \frac{4\pi\alpha}{f_V{}^2} g_V(E_\gamma)\sigma_t(VN)\,, \tag{4.12}$$

where

$$g_V(E_\gamma)=\frac{m_V{}^4}{4k_V^{02}(k_V^0-k^0)^2}\,\frac{p^0+k_V^0}{p^0+k^0}\,. \tag{4.13}$$

Equation (4.12) gives, in the limit of $E_\gamma\to\infty$,

$$\sigma_t(\gamma N)|_{E_\gamma\to\infty}=99\ \mu\text{b}\,, \tag{4.14}$$

which may be compared with the experimental value[29] 97.4 ± 1.9 μb. Obviously in (4.14) we do not have any parameters to adjust; these have been determined in the last section. If we take the Compton scattering from the nucleons as almost wholly diffractive,[30] we then obtain from (4.14) in the asymptotic limit

$$\sigma_t^{\text{el}}(\gamma N)=\frac{\sigma_t(\gamma N)^2}{16\pi}\,\frac{1}{b}\simeq 66\ \text{nb}\,, \tag{4.15}$$

where we have taken $b=7.6$ GeV^{-2}, as the ρ^0 contribution is dominant. The energy dependence of this process is as given by (4.13), and is not good at low energies since the maximum contribution comes from ρ^0 and ω, and for these processes until moderate energies, other processes which are not diffractive continue to contribute,[31] as was seen in Sec. IIIA. It is not possible to do any better without a determination of the correct energy dependence for vector mesons at low energies, a problem we have not tackled in this paper. Without details regarding this, and without considering the process (4.3), we have to remain satisfied with the asymptotic result (4.14).

V. DISCUSSIONS

We note that it has been possible to understand the vector-dominance model for electromagnetic interactions as a reflection of the fact that the contribution from the direct photon absorption by a quark in a hadron is very small at high energies. This is true for the process $\gamma p\to Vp$, where one quark will absorb the photon and there will be two spectator quarks, giving rise to a small overlap integral from the wave functions of the hadrons.[5] On the other hand, contribution through the exchange of a vector meson continues to be large, resulting essentially in a vector-meson-dominance model with kinematic corrections. The problem here has *not* been resolved at moderate energies (say up to $E_\gamma=4$ GeV), since we have concentrated our attention on *only* the process with a vector-meson exchange. Even within the scope of this model we also believe that corrections including mixings may be useful if we want a more complete fit of the data at the above energies.[32] The vector-dominance model here arises from the quark model through nonvanishing of the matrix element $\langle\text{vac}|J^\mu(0)|V\lambda\rangle$. This model is different from conventional vector-dominance models since the photon interacts with hadrons at the quark level directly, as illustrated in Ref. 5 at low energies, whereas at high energies vector dominance comes into play. In attempting this, we observe that the modified model automatically generates an energy dependence of $(d\sigma/dt)_0$ and slope parameters for photoproduction of vector mesons which are in good agreement with experiments.

We note that we have not considered the excited levels of the vector mesons. In (4.9) we have explicitly calculated the effects of successive vector mesons and find that the contributions become extremely small quite fast. When one includes the corrections due to mixing,[32] it may be worthwhile to take this into account. However, at the present level one has merely found a good reason for considering the vector-dominance model as a very reasonable guess in the context of the quark model[5] as stated above.

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