[6]In all our coherent calculations we assume a simple exponential charge distribution rather than the possibly more realistic Fermi distribution to calculate the form factor. As a result, our coherent cross sections are too large. Near threshold our results are too large by an order of magnitude and well above threshold our results are about 20% larger than those done assuming a Fermi charge distribution. Near threshold the coherent cross section is a negligible part of the total cross section and thus is unimportant. In general, because of the severe κ dependence of the cross sections, choice of coherent form factors is a matter of preference.

[7]R. Linsker, Columbia University Report No. NYO-1932(2)-199, 1971 (unpublished).

[8]J. Reiff, Nucl. Phys. B28, 495 (1971).

[9]The data for W_1-production cross sections are taken from R. W. Brown and J. Smith [Phys. Rev. D 3, 207 (1971)] and also from calculations done by P. J. Oddone, in NAL Report No. 21 (unpublished).

[10]The value of $M_1 = 37.29$ GeV/c^2 was included since T. D. Lee has recently proposed this value for the W_1 mass [T. D. Lee, Phys. Rev. Letters 26, 801 (1971)].

[11]The data here for W_1 production used in comparison with W_0 production are taken from R. W. Brown and J. Smith, Ref. 9, and also from P. J. Oddone (private communication).

[12]These data for W_1 production used in comparison with W_0 production are from R. W. Brown, R. H. Hobbs, and J. Smith, Phys. Rev. D 4, 794 (1971).

[13]For a more quantitative explanation, see J. S. Bell and M. Veltman, Phys. Letters 5, 151 (1963).

[14]The q^2 in this section should not be confused with the electromagnetic $q^2 = t$ mentioned earlier. W production can be considered from the point of view of neutrino inelastic scattering, and then q^2 is the momentum transfer squared at the lepton vertex, which is equal to $|(\nu - \mu)^2|$.

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[16]The results presented in this paper are for elastic and coherent production only. The inelastic process will increase the total cross section, but should be characterized by the same qualitative behavior in its differential distributions. This follows since the distributions depend most critically upon which diagram dominates and not the details of the hadronic electromagnetic vertex.

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Search for New Vector Mesons by Diffractive Photoproduction at the National Accelerator Laboratory

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If an as yet unobserved vector meson exists, it can be seen most easily by diffractive photoproduction on nuclei using a photon beam at the National Accelerator Laboratory (NAL). We have considered the problem of coherent production of such a particle on nuclei, taking into account the mixing between the ρ and the new vector meson to arbitrary order. The method used here can be generalized to describe coherent scattering of hadrons and nuclei. It will be shown that the interference effect reduces the production of the new vector meson considerably and that the A dependence of the production cross section plays an important role.

INTRODUCTION

In order to formulate the concept of vector-meson dominance, it is crucial to know whether or not there are other vector mesons besides the ρ, ω, and ϕ through which the photon interacts with hadrons. Since ρ, ω, and ϕ are produced diffractively when a high-energy photon beam strikes a nucleus, we expect such mesons to also be photoproduced diffractively. The diffractive production has the advantage that the produced particle must have quantum numbers similar to those of the pho-

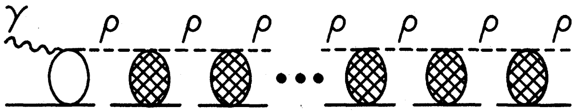

FIG. 1. The diagram for ρ production. The solid line corresponds to various nucleons.

ton and, furthermore, the nondiffractive background is orders of magnitude smaller in the forward direction. We shall call the new vector meson v. An attempt to photoproduce v was made and no such meson was found with its mass $m_v \leq 2$ GeV.[1,2] At presently available energies, 2 GeV is an upper limit on the mass for diffractive production.

Since a photon beam with energy as high as 300 GeV will be available at the National Accelerator Laboratory (NAL), questions on the existence of v with $m_v \lesssim 7$ GeV will be settled soon. Some theoretical investigations must be made before such an experiment to look for v is planned, since there are effects which may greatly reduce chances of detecting v even if it exists. For example:

(i) Absorption effect. Since v is a strong-interacting particle, the final-state interaction will decrease the probability of observing v. Such an effect is seen in ρ photoproduction.[3]

(ii) Interference effect. Since ρ and v have the same quantum numbers, these two states will mix with each other.[4] This effect may be such that the probability for observing v is greatly reduced.

Since the dynamics of v is unknown, we shall not be able to give an absolute statement on these possibilities. We are able, however, to express the cross section for

$$\gamma + \text{nucleus} \to v + \text{nucleus}$$

in terms of four unknown parameters m_v, f_v,

$$\sigma_v = v + p \text{ total cross section},$$

and

$$\frac{ik}{4\pi}\Sigma = v + p \to \rho + p \text{ forward scattering amplitude},$$

where p stands for a proton and k is the incident momentum of the photon. We expect Σ to be real and constant at high energy. The results of the calculation indicate that for any value of Σ the interference effect reduces the diffractive production of v considerably and, depending on the value of σ_v, the absorption effect will be appreciable. It also indicates, however, that if v existed with $m_v \lesssim 7$ GeV, we should be able to observe them at NAL.[5]

I. FORMULATION

The theory of photoproduction on a nucleus is a many-body problem and is very complex depending

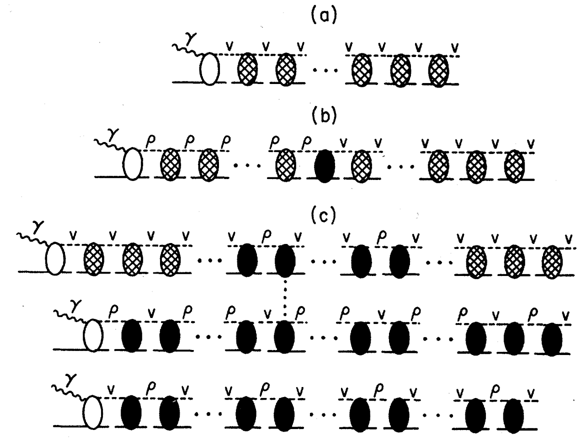

FIG. 2. (a) A diagram for the v production corresponding to the ρ production of Fig. 1. The cross-hatched area corresponds to elastic diffractive scattering. (b) A diagram for the v production including one $\rho \sim v$ transition. The solid area corresponds to diffractive $\rho p \rightleftarrows vp$ scattering. (c) Set of diagrams for the v including arbitrary number of ρ-v transitions.

on the detail required. Therefore, the main problem is to realize the limitation placed on the experimental measurement which is to be compared with the theory and reduce the complexity of the problem by a well-chosen series of approximations. The general treatment of the scattering theory at high energy is given by Glauber,[6] and subsequent application to ρ photoproduction on nuclei neglecting all correlation of nucleons has been shown to be quite adequate for explaining present experimental data. The basic set of assumptions successfully used in ρ production, therefore, should also be applicable in v photoproduction.

Let us consider the process of ρ photoproduction in the absence of v and compare the problem with that of v photoproduction. In Fig. 1 we have drawn diagrams for ρ production. The time runs from left to right. The solid line corresponds to various nucleons in the nucleus. Note that the diffractive production implies that a majority of scattering is forward, $\theta \sim 1/ka$ being the width of the forward peak, where a is the effective radius of a nucleon and k is the momentum of the photon beam. Taking k to be 200 GeV/c and a to be 1 F, we have $1/ka \sim \frac{1}{40}$. So, to a good approximation, we can take each scattering of ρ on nucleon to be forward. Thus, for each impact parameter of the photon, we have a one-dimensional problem. To show this, nucleons are drawn in a straight line. Since the backward direction is negligible, each nucleon participates in the scattering once. The break in the solid line indicates that they are different nucleons.

Now consider the v photoproduction. The photon creates either ρ or v (we denote it by α) at $\vec{R} = (\vec{b}, z)$ by scattering diffractively with one of the

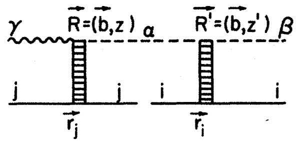

FIG. 3. The photon enters a nucleus and scatters diffractively with the jth nucleon at $\vec{r}_j$ to create α at $\vec{R}$. α in turn scatters diffractively with the ith nucleon at $\vec{r}_i$ and creates β at $\vec{R}'$. The ladder corresponds to interaction.

nucleons, where $\vec{b}$ denotes the impact parameter and the z axis is in the direction of the beam. To lowest order in $e^2/4\pi$, the photon no longer enters into the problem. Then α scatters diffractively with each of the other nucleons in the nucleus, but we must consider the possibility that v and ρ are mixed arbitrarily many times. The processes mentioned here are shown in Fig. 2.

For simplicity let us consider a problem in which the photon scatters diffractively with the nucleon j at $\vec{r}_j$; α is created at $\vec{R}=(\vec{b},z)$ and in turn scatters with nucleon i at $\vec{r}_i$ and turns into β, where α and β again stand for either ρ or v. This process is shown in Fig. 3. Let $f_{\alpha\beta}(\vec{q})$ denote the scattering amplitude for $\alpha+p\to\beta+p$ with momentum transfer $\vec{q}$ and $\Gamma_{\alpha\beta}(\vec{r})$ be its Fourier transform:

$$\Gamma_{\alpha\beta}(\vec{r})=\frac{1}{(2\pi)^{3/2}ik}\int e^{-i\vec{q}\cdot\vec{r}}f_{\alpha\beta}(\vec{q})\,d^3q, \tag{1}$$

$$f_{\alpha\beta}(\vec{q})=\frac{ik}{(2\pi)^{3/2}}\int e^{i\vec{q}\cdot\vec{r}}\Gamma_{\alpha\beta}(\vec{r})\,d^3r. \tag{2}$$

The scattering amplitude for the event is

$$\begin{aligned}f^{ji}_{\gamma\beta}(\vec{Q})=\frac{ik}{(2\pi)^{3/2}}\int d^3R\prod_{k=1}^{A}d^3r_k\\ \times e^{i\vec{Q}\cdot\vec{R}}\,U(\vec{r}_1,\ldots,\vec{r}_A)^*\Gamma_{\gamma\alpha}(\vec{R}-\vec{r}_j)\\ \times\exp[i\chi^i_{\alpha\beta}(\vec{b},z,q^{\parallel}_{\alpha\beta})]\,U(\vec{r}_1,\ldots,\vec{r}_A),\end{aligned} \tag{3}$$

where $\vec{Q}$ is the momentum transfer between γ and β, $\chi^i_{\alpha\beta}(\vec{b},z,q^{\parallel}_{\alpha\beta})$ is the phase function describing $\alpha+i$th nucleon $\to\beta+i$th nucleon scattering. $q^{\parallel}_{\alpha\beta}$ is the longitudinal momentum transfer

$$[\vec{q}_{\alpha\beta}]_{\text{forward direction}}=\frac{m_\beta^2-m_\alpha^2}{2k}\hat{z}.$$

The state of the nucleus is $U(\vec{r}_1,\ldots,\vec{r}_A)$. To obtain Eq. (3), first note that $\Gamma_{\gamma\alpha}(\vec{R}-\vec{r}_j)\times\exp[i\chi^i_{\alpha\beta}(\vec{b},z,q^{\parallel}_{\alpha\beta})]$ describes the double scattering $\gamma+p\to\alpha+p$ and $\alpha+p\to\beta+p$ as a function of $\vec{r}_j$, $\vec{r}_i$, and $\vec{R}$. We then sandwich the above expression between the states of the nucleus, integrate over all the internal nuclear coordinates $\vec{r}_k$, and finally take the Fourier transform to obtain the scattering amplitude. To find $\chi^i_{\alpha\beta}(\vec{b},z,q^{\parallel}_{\alpha\beta})$, we write the $\alpha+i$th nucleon $\to\beta+i$th nucleon scattering amplitude

$$\begin{aligned}f^i_{\alpha\beta}(\vec{q})=\frac{ik}{(2\pi)^{3/2}}\int e^{i\vec{q}\cdot\vec{b}}\,d^2b\,e^{iq^{\parallel}_{\alpha\beta}z'}\,U(\vec{r}_1,\ldots,\vec{r}_A)^*\\ \times\Gamma_{\alpha\beta}(\vec{b}-\vec{b}_i,z'-z_i)\\ \times U(\vec{r}_1,\ldots,\vec{r}_A)dz'\prod_{k=1}^{A}d^3r_k.\end{aligned} \tag{4}$$

In terms of the phase $\chi^i_{\alpha\beta}(\vec{b}-\vec{b}_i,q^{\parallel}_{\alpha\beta})$, the same amplitude can be written as

$$\begin{aligned}f^i_{\alpha\beta}(\vec{q})=\frac{ik}{2\pi}\int d^2b\,e^{i\vec{q}\cdot\vec{b}}\,U(\vec{r}_1,\ldots,\vec{r}_A)^*\\ \times\{\delta_{\alpha\beta}-\exp[i\chi^i_{\alpha\beta}(\vec{b}-\vec{b}_i,q^{\parallel}_{\alpha\beta})]\}\\ \times U(\vec{r}_1,\ldots,\vec{r}_A)\prod_{k=1}^{A}d^3r_k.\end{aligned} \tag{5}$$

Comparing Eqs. (4) and (5), we obtain

$$\begin{aligned}\exp[i\chi^i_{\alpha\beta}(\vec{b}-\vec{b}_i,q^{\parallel}_{\alpha\beta})]\\ =\delta_{\alpha\beta}-\frac{1}{(2\pi)^{1/2}}\int_{-\infty}^{\infty}e^{iq^{\parallel}_{\alpha\beta}z'}\Gamma_{\alpha\beta}(\vec{b}-\vec{b}_i,z'-z_i)\,dz'.\end{aligned} \tag{6}$$

The $\chi^i_{\alpha\beta}(\vec{b},q^{\parallel}_{\alpha\beta})$ of Eq. (6) is the phase for the process in which α is present at the entire range of z', $-\infty<z'<\infty$. But, in fact, in our process α is created at $\vec{R}=(\vec{b},z)$, so z' runs in the range $z<z'<\infty$ and we must cut off the z' integration at z. Then Eq. (3) becomes

$$\begin{aligned}f^{ji}_{\gamma\beta}(\vec{Q})=\frac{ik}{(2\pi)^{3/2}}\int d^3R\,e^{i\vec{Q}\cdot\vec{R}}\,U(\vec{r}_1,\ldots,\vec{r}_A)^*\Gamma_{\gamma\alpha}(\vec{R}-\vec{r}_j)\\ \times\left(\delta_{\alpha\beta}-\frac{1}{(2\pi)^{1/2}}\int_z^{\infty}e^{iq^{\parallel}_{\alpha\beta}z'}\Gamma_{\alpha\beta}(\vec{b}-\vec{b}_i,z'-z_i)\,dz'\right)U(\vec{r}_1,\ldots,\vec{r}_A)\prod_{k=1}^{A}d^3r_k.\end{aligned} \tag{7}$$

Using the fact that

$$\int U(\vec{r}_1, \ldots, \vec{r}_A)^* U(\vec{r}_1, \ldots, \vec{r}_A) \prod_{k=1}^{A} d^3 r_k = 1,$$

$$\int U(\vec{r}_1, \ldots, \vec{r}_A)^* U(\vec{r}_1, \ldots, \vec{r}_A) \prod_{\substack{k=1 \\ k\neq i}}^{A} d^3 r_k = \rho_i(\vec{r}_i), \tag{8}$$

and

$$\int U(\vec{r}_1, \ldots, \vec{r}_A)^* U(\vec{r}_1, \ldots, \vec{r}_A) \prod_{\substack{k=1 \\ k\neq i, k\neq j}}^{A} d^3 r_k = \rho_{ij}(\vec{r}_i, \vec{r}_j),$$

where $\rho_i(\vec{r}_i)$ is the distribution function for the ith nucleon and $\rho_{ij}(\vec{r}_i, \vec{r}_j)$ is the distribution function for the jth and ith nucleons, we can rewrite Eq. (7) as

$$f^{ji}_{\gamma\beta}(\vec{Q}) = \frac{ik}{(2\pi)^{3/2}} \int d^3R\, d^3r_j\, e^{i\vec{Q}\cdot\vec{R}}\, \Gamma_{\gamma\alpha}(\vec{R} - \vec{r}_j) \times \left(\delta_{\alpha\beta}\rho_j(\vec{r}_j) - \frac{1}{(2\pi)^{1/2}} \int_z^\infty \Gamma_{\alpha\beta}(\vec{b} - \vec{b}_i, z' - z_i) \times e^{iq^{\parallel}_{\alpha\beta} z'} dz'\, \rho_{ij}(\vec{r}_i, \vec{r}_j)\, d^3r_i \right), \tag{9}$$

where we have integrated over the spectator nucleons, $k = 1, \ldots, A$, $k \neq i, j$. We now make an approximation that $\rho_j(\vec{r}_j) = \rho(\vec{r}_j)$ for all j. We have neglected all correlation effects between nucleons. The correlation effects may not be negligible. The accuracy of the experimental measurement, however, does not seem to require nuclear-physics considerations at least for low-energy ρ photoproduction, and therefore we will adopt this approximation. The $\rho(\vec{r}_i)$ is equal to the average density of the nucleus and

$$U(\vec{r}_1, \ldots, \vec{r}_A)^* U(\vec{r}_1, \ldots, \vec{r}_A) = \prod_{k=1}^{A} \rho(\vec{r}_k).$$

Since $\Gamma_{\alpha\beta}(\vec{b} - \vec{b}_i, z' - z_i) \neq 0$ only in the immediate neighborhood of the ith nucleon, for large A, $\rho(\vec{r}_i)$ varies little over the same region. Then we may make an approximation

$$\int d^3x\, d^3y\, e^{i\vec{k}\cdot\vec{x}}\, \Gamma_{\alpha\beta}(\vec{x} - \vec{y}) \rho(\vec{y}) = \int d^3x\, e^{i\vec{k}\cdot\vec{x}} \rho(\vec{x}) \int d^3y\, \Gamma_{\alpha\beta}(\vec{x} - \vec{y}) = \frac{(2\pi)^{3/2}}{ik} f_{\alpha\beta}(0) \int d^3x\, e^{i\vec{k}\cdot\vec{x}} \rho(\vec{x}).$$

Using this, we can rewrite Eq. (9) as

$$f^{ji}_{\gamma\beta}(\vec{Q}) = f_{\gamma\alpha}(0) \int d^3R\, e^{i\vec{Q}\cdot\vec{R}} \rho(\vec{R}) \times \left(\delta_{\alpha\beta} - \frac{2\pi}{ik} f_{\alpha\beta}(0) \int_z^\infty dz'\, e^{iq^{\parallel}_{\alpha\beta} z'} \rho(\vec{b}, z') \right). \tag{10}$$

So far we have been considering a hypothetical problem in which we have turned off the interaction of α with all nucleons except the ith nucleon. The 2×2 matrix

$$T^i_{\alpha\beta}(\vec{R}) = \begin{bmatrix} 1 - \frac{2\pi}{ik} f_{\rho\rho}(0) \int_z^\infty \rho(\vec{b}, z')\, dz' & -\frac{2\pi}{ik} f_{\rho v}(0) \int_z^\infty e^{iq^{\parallel}_{\rho v} z'} \rho(\vec{b}, z')\, dz' \\ -\frac{2\pi}{ik} f_{\rho v}(0) \int_z^\infty e^{-iq^{\parallel}_{\rho v} z'} \rho(\vec{b}, z')\, dz' & 1 - \frac{2\pi}{ik} f_{vv}(0) \int_z^\infty \rho(\vec{b}, z')\, dz' \end{bmatrix} \tag{11}$$

describes the scattering between α and the ith nucleon. Note that $T^i_{\alpha\beta}(\vec{R})$ does not depend on i. That is, it does not depend on which nucleon α scatters once the nuclear correlation is neglected. The time ordering for interaction with different nucleons does not occur, since interaction operations all commute with each other. We have $f_{\rho v}(0) = f_{v\rho}(0)$ using time-reversal invariance. Furthermore, if we assume that all $f_{\alpha\beta}(0)$ become purely imaginary at high energy, $T^i_{\alpha\beta}(\vec{R})$ is Hermitian. Then, when we include the effect of α scattering with all other nucleons, $T_{\alpha\beta}(\vec{R})$ is replaced by $[T(\vec{R})^{A-1}]_{\alpha\beta}$. Let $\chi^{\pm}(\vec{R})$ be the eigenvector for $T(\vec{R})$ with eigenvalue $\lambda^{\pm}(\vec{R})$. Denoting

$$T(\vec{R}) = \begin{pmatrix} a & b \\ b^* & d \end{pmatrix}, \quad \lambda^{\pm} = \tfrac{1}{2}\{a + d \pm [(a-d)^2 + 4|b|^2]^{1/2}\},$$

and in terms of eigenvectors $|\chi^{\pm}\rangle$, ρ and v can be written as

$$|\rho\rangle = \frac{(a-\lambda^-)[|b|^2 + (a-\lambda^+)^2]^{1/2}}{(\lambda^+ - \lambda^-)b} |\chi^+\rangle + \frac{(a-\lambda^+)[|b|^2 + (a-\lambda^-)^2]^{1/2}}{(\lambda^+ - \lambda^-)b} |\chi^-\rangle,$$
$$|v\rangle = \frac{[|b|^2 + (a-\lambda^+)^2]^{1/2}}{\lambda^+ - \lambda^-} |\chi^+\rangle + \frac{[|b|^2 + (a-\lambda^-)^2]^{1/2}}{\lambda^+ - \lambda^-} |\chi^-\rangle. \tag{12}$$

When the state $\chi^{\pm}(\vec{R})$ interacts with a nucleon it becomes $\lambda^{\pm}(\vec{R})\chi^{\pm}(\vec{R})$. The influence of $A-1$ nucleons will be $[\lambda^{\pm}(\vec{R})]^{A-1}\chi^{\pm}(\vec{R})$. Thus, the effects of all the nucleons can be treated in a trivial way when the nuclear-correlation effect is neglected.

The amplitude for photoproduction of β on a nucleus with momentum transfer $\vec{Q}$ can be obtained by generalizing Eq. (10):

$$f_{\gamma A\to\beta A}(\vec{Q})=\sum_{\alpha=\rho,v}\sum_{j=1}^{A}f_{\gamma\alpha}(0)\int d^3R\, e^{i\vec{Q}\cdot\vec{R}}\rho(\vec{R})[T(\vec{R})^{A-1}]_{\alpha\beta}$$
$$=2\pi A\sum_{\alpha=\rho,v}f_{\gamma\alpha}(0)\int e^{iQ_\beta^{\parallel}z}J_0(kb\theta)\rho(\vec{b},z)[T(\vec{R})^{A-1}]_{\alpha\beta}b\,db\,dz\,, \tag{13}$$

where $Q_\beta^{\parallel}=m_\beta{}^2/2k$ and θ is the angle that β makes with respect to the incident photon beam. The second expression is obtained by integrating over the azimuthal angle taking advantage of the cylindrical symmetry.[7] Denoting

$$|\rho\rangle=E|\chi^+\rangle+B|\chi^-\rangle,$$
$$|v\rangle=C|\chi^+\rangle+D|\chi^-\rangle, \tag{14}$$

where the coefficients stand for those given in Eq. (12), we obtain expressions for photoproduction of ρ and v on a nucleus from Eq. (13):

$$f_{\gamma A\to vA}(\theta)=2\pi A\int b\,db\,dz\,e^{iq_v^{\parallel}z}\rho(\vec{b},z)J_0(kb\theta)\{f_{\gamma\rho}(0)[(\lambda^+)^{A-1}E^*C+(\lambda^-)^{A-1}B^*D]+f_{\gamma v}(0)[(\lambda^+)^{A-1}C^*C+(\lambda^-)^{A-1}D^*D]\}, \tag{15}$$

$$f_{\gamma A\to\rho A}(\theta)=2\pi A\int b\,db\,dz\,e^{iq_\rho^{\parallel}z}\rho(\vec{b},z)J_0(kb\theta)\{f_{\gamma\rho}(0)[(\lambda^+)^{A-1}|E|^2+(\lambda^-)^{A-1}|B|^2]+f_{\gamma v}(0)[(\lambda^+)^{A-1}C^*E+(\lambda^-)^{A-1}D^*B]\}. \tag{16}$$

A similar technique can be used to derive the expression

$$f_{\rho A\to vA}(\theta)=\frac{k}{i}\int b\,db\,J_0(kb\theta)[(\lambda^+)^A\tilde{E}^*\tilde{C}+(\lambda^-)^A\tilde{B}^*\tilde{D}], \tag{17}$$

where $\tilde{E}$, $\tilde{B}$, $\tilde{C}$, and $\tilde{D}$ are calculated in a similar way as E, B, C, and D, respectively, with the exception that the lower limit of the integral in the matrix element of Eq. (11), z, is replaced by $-\infty$.

II. PARAMETERS

Equations (15) and (16) contain parameters

$$f_{\gamma\rho}(0),\, f_{\gamma v}(0),\, f_{\rho v}(0)=f_{v\rho}(0),\, f_{\rho\rho}(0),\, f_{vv}(0),$$

where $f_{\alpha\beta}(0)$ is the $\alpha+p\to\beta+p$ scattering amplitude at zero momentum transfer. For $\alpha\neq\beta$, $f_{\alpha\beta}(0)$ is an unphysical amplitude which is obtained by extrapolation from the physical region where $|\vec{q}|\geq|(m_\beta{}^2-m_\alpha{}^2)/2k|$. $f_{\gamma\rho}(0)$ and $f_{\gamma v}(0)$ can be written in terms of $f_{\rho\rho}(0)$, $f_{vv}(0)$, and $f_{\rho v}(0)$ using the vector-dominance model as shown in Fig. 4:

$$f_{\gamma\rho}(0)=\frac{e}{f_\rho}f_{\rho\rho}(0)+\frac{e}{f_v}f_{v\rho}(0),$$
$$f_{\gamma v}(0)=\frac{e}{f_\rho}f_{\rho v}(0)+\frac{e}{f_v}f_{vv}(0)\,. \tag{18}$$

Since existence of v is not required by the usual vector-dominance model at present energies, we expect $1/f_\rho\gg 1/f_v$.

The simplest kind of diffraction theory (scattering from a black disk or simple Pomeranchuk-Regge-pole exchange) predicts that $f_{\alpha\beta}(0)$ is purely imaginary at high energy. ρ-photoproduction data seem to require $\mathrm{Re}f_{\gamma\rho}(0)/\mathrm{Im}f_{\gamma\rho}(0)\approx 0.2$ at $k\lesssim 15$ GeV/c. If we are interested only in an order-of-magnitude prediction for the v production, we can set the amplitude to be purely imaginary. Thus, using the optical theorem, we have

$$\frac{4\pi}{ik}f_{\rho\rho}(0)=\sigma_\rho,\quad \frac{4\pi}{ik}f_{vv}(0)=\sigma_v\,,$$

and define

$$\frac{4\pi}{ik}f_{\rho v}(0)\equiv\Sigma\,, \tag{19}$$

where σ_ρ and σ_v are ρp and vp total cross sections, respectively. We take σ_ρ, σ_v, and Σ to be independent of energy. This is consistent with the approximation of setting the amplitude to be purely imaginary. From ρ photoproduction[8] we have $\sigma_\rho\approx 26$ mb. Therefore, we have σ_v, Σ, $f_v(0)$, and m_v as unknown parameters.

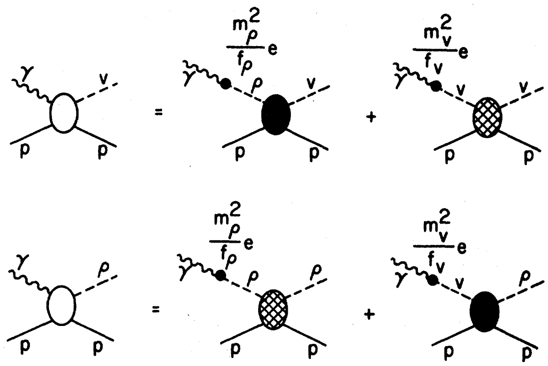

FIG. 4. Diagrams for $f_{\gamma v}$ and $f_{\gamma\rho}$.

III. SMALL-Σ AND LARGE-k LIMIT

Although we cannot *a priori* restrict ourselves to small Σ and large k, by studying these limits we gain insight into our general result. In this section we intend to show that (a) if the absorption of the vector mesons due to scattering with nucleons is very small, i.e., σ_ρ and σ_v are small, the A dependence of the cross section for photoproduction of both ρ and v goes as A^2, and (b) if the absorption is large, with ρ dominance $1/f_v=0$, the A dependence of the cross section for the ρ production is $A^{4/3}$, while the A dependence of the v-production cross section may be quite different, depending on the value for σ_v. On the other hand, if $1/f_v \neq 0$, the v-production cross section will have a term which behaves as $A^{4/3}$. We expect, therefore, that a study of the A dependence of v production will be a very sensitive way to measure the parameters f_v and Σ.

As $\Sigma \to 0$, $b \to 0$. To first order in b, we have, from Eq. (12), $\lambda^+ \cong a$, $\lambda^- \cong d$,

$$E^*C=\frac{b}{a-d}\,,\quad B^*D=\frac{-b}{a-d}\,,\quad |C|^2\cong 0\,,\quad |D|^2=1\,.$$

Then

$$(\lambda^+)^{A-1}=\left(1-\tfrac{1}{2}\sigma_\rho\int_z^\infty \rho(\vec{b},z)\,dz\right)^{A-1}$$
$$\cong \exp\left(-\tfrac{1}{2}\sigma_\rho A\int_z^\infty \rho(\vec{b},z')\,dz'\right).$$

Substituting these into Eq. (15), we obtain

$$f_{\gamma A\to vA}(\theta)=2\pi A\int b\,db\,dz\,e^{iq_v^{\parallel}z}\rho(\vec{b},z)J_0(kb\theta)\Bigg\{f_{\gamma\rho}(0)\,\frac{-\Sigma\int_z^\infty e^{iq_v^{\parallel}z'}\rho(\vec{b},z')\,dz'}{(\sigma_v-\sigma_\rho)\int_z^\infty\rho(\vec{b},z')\,dz'}$$
$$\times\left[\exp\left(-\tfrac{1}{2}\sigma_\rho A\int_z^\infty\rho(\vec{b},z')\,dz'\right)-\exp\left(-\tfrac{1}{2}\sigma_v A\int_z^\infty\rho(\vec{b},z')\,dz'\right)\right]+f_{\gamma v}(0)\exp\left(-\tfrac{1}{2}\sigma_v A\int_z^\infty\rho(\vec{b},z')\,dz'\right)\Bigg\}\,. \tag{20}$$

In the limit as $k\to\infty$ and $q_v^{\parallel}=0$, Eq. (20) simplifies considerably. The z integration can be performed, and we obtain

$$f_{\gamma A\to vA}(\theta)=4\pi\int b\,db\,J_0(kb\theta)\Bigg\{\frac{f_{\gamma\rho}(0)\Sigma}{\sigma_\rho-\sigma_v}\left(\frac{1-\exp\left(-\tfrac{1}{2}\sigma_\rho A\int_{-\infty}^\infty\rho(\vec{b},z')\,dz'\right)}{\sigma_\rho}-\frac{1-\exp\left(-\tfrac{1}{2}\sigma_v A\int_{-\infty}^\infty\rho(\vec{b},z')\,dz'\right)}{\sigma_v}\right)$$
$$+\frac{f_{\gamma v}(0)}{\sigma_v}\left[1-\exp\left(-\tfrac{1}{2}\sigma_v A\int_{-\infty}^\infty\rho(\vec{b},z')\,dz'\right)\right]\Bigg\}\,. \tag{15$'$}$$

Similarly in this limit, Eq. (16) becomes

$$f_{\gamma A\to\rho A}=\frac{4\pi}{\sigma_\rho}f_{\gamma\rho}(0)\int b\,db\,J_0(kb\theta)\left[1-\exp\left(-\tfrac{1}{2}\sigma_\rho A\int_{-\infty}^\infty\rho(\vec{b},z')\,dz'\right)\right] \tag{16$'$}$$

and Eq. (17) becomes

$$f_{\rho A\to vA}=\frac{-4\pi f_{\rho v}(0)}{\sigma_v-\sigma_\rho}\int b\,db\,J_0(kb\theta)\left[\exp\left(-\tfrac{1}{2}\sigma_v A\int_{-\infty}^\infty\rho(\vec{b},z')\,dz'\right)-\exp\left(-\tfrac{1}{2}\sigma_\rho A\int_{-\infty}^\infty\rho(\vec{b},z')\,dz'\right)\right]. \tag{17$'$}$$

Equation (17′) is the result obtained previously by other authors.[9] In Eq. (17′), replace ρ by γ and v by ρ; then setting $\sigma_\gamma=O(\alpha^2)\ll\Sigma,\sigma_\rho$, we obtain Eq. (16′). By comparing the second term of Eq. (15′) with Eq. (16′) we see that this term is direct photoproduction of v. Diagrammatically, it is shown in Fig. 2(a). It can also be shown that the first term of Eq. (15′) corresponds to the diagram shown in Fig. 2(b). Therefore, the $\Sigma\to 0$ limit given in Eq. (15′) corresponds to the first-order perturbation expansion in Σ. Finally, under the ρ-dominance assumption $1/f_v=0$, we expect

$$f_{\gamma A\to vA}(\theta)=(1/f_\rho)f_{\rho A\to vA}(\theta)$$

to hold. This is indeed satisfied, as shown below.

(a) Consider the limit of small σ_ρ and σ_v. This limit corresponds to the nucleus being transparent to ρ and v. Taking $\theta=0$, expanding the exponentials, and performing the integration over b, we obtain

$$\begin{aligned} f_{\gamma A\to vA}(0)&=Af_{\gamma v}(0)\,,\\ f_{\gamma A\to\rho A}(0)&=Af_{\gamma\rho}(0)\,,\\ f_{\rho A\to vA}(0)&=Af_{\rho v}(0)\,. \end{aligned}\tag{21}$$

The cross section, therefore, goes as A^2 as expected from the coherent scattering.

(b) Consider large σ_ρ and σ_v. This limit corresponds to the nucleus being black. Since the exponential is negligible compared to unity in Eq. (16′), it becomes

$$f_{\gamma A\to\rho A}(0)\to\frac{1}{\sigma_\rho}f_{\gamma\rho}2\pi R^2=\frac{e}{f_\rho}\,\frac{ik}{4\pi}\,2\pi R^2\tag{22}$$

to the zeroth order in Σ.

Since $2\pi R^2$ is the total cross section for diffractive scattering from a black disk, using the optical theorem, it is consistent with $\rho A\to\rho A$ diffraction scattering from a black disk. In order to understand this result in terms of well-known diffraction concepts consider

$$f_{\alpha A\to\beta A}(q)=\frac{ik}{2\pi}\int e^{i\vec{q}\cdot\vec{b}}(\delta_{\alpha\beta}-e^{i\chi_{\alpha\beta}(\vec{b})})d^2b\,.\tag{23}$$

Then the phase function in this limit is

$$\begin{aligned} e^{i\chi_{\rho\rho}(\vec{b})}&=1\quad\text{for } |\vec{b}|>R\\ &=0\quad\text{for } |\vec{b}|<R\,. \end{aligned}\tag{24}$$

The limit for Eq. (17′) is $f_{\rho A\to vA}(0)\to 0$. In terms of Eq. (23), the phase function in this limit is

$$e^{i\chi_{\rho v}(\vec{b})}=0\quad\text{for all }\vec{b}\,.$$

This is because for $|\vec{b}|>R$ there is no mechanism which transforms ρ to v, and for $|\vec{b}|<R$ all v created are absorbed in the nucleus and are never detected.

The behavior of Eq. (15′) is our main interest. In this limit, the first term of Eq. (15′) becomes

$$-2\pi f_{\gamma\rho}(0)\,\frac{\Sigma}{\sigma_\rho-\sigma_v}\left(\frac{1}{\sigma_v}-\frac{1}{\sigma_\rho}\right)R^2=-\frac{ike}{2}\,\frac{1}{f_\rho}\,\frac{\Sigma}{\sigma_v}R^2$$

and the second term of Eq. (15′) becomes

$$2\pi\,\frac{f_{\gamma v}(0)}{\sigma_v}\,R^2=\frac{ike}{2}\left(\frac{1}{f_\rho}\,\frac{\Sigma}{\sigma_v}\right)R^2\,,$$

assuming $1/f_v=0$. If we assume ρ dominance, $1/f_v=0$, the first and second terms cancel each other. That is, in the lowest order in Σ, the diagrams shown in Figs. 2(a) and 2(b) interfere destructively to cancel out the $A^{4/3}$ behavior. Substituting Eq. (18) into Eq. (15′), we obtain

$$\begin{aligned} f_{\gamma A\to vA}(\theta)=\;&\frac{e}{f_\rho}\,\frac{4\pi f_{\rho v}(0)}{\sigma_v-\sigma_\rho}\int b\,db\,J_0(kb\theta)\left[\exp\left(-\tfrac{1}{2}\sigma_\rho A\int_{-\infty}^{\infty}\rho(\vec{b},z')\,dz'\right)-\exp\left(-\tfrac{1}{2}\sigma_v A\int_{-\infty}^{\infty}\rho(\vec{b},z')\,dz'\right)\right]\\ &+\frac{e}{f_v}\,\frac{4\pi f_{vv}(0)}{\sigma_v}\int b\,db\,J_0(kb\theta)\left[1-\exp\left(-\tfrac{1}{2}\sigma_v A\int_{-\infty}^{\infty}\rho(\vec{b},z')\,dz'\right)\right]. \end{aligned}\tag{25}$$

If $1/f_v=0$, comparing Eq. (25) with Eq. (17′), we obtain

$$f_{\gamma A\to vA}(0)=(e/f_\rho)f_{\rho A\to vA}\,,$$

which is consistent with ρ dominance.

Let us examine the A dependence of Eq. (25) in some detail. For this purpose, it is convenient to take a cylindrical model for the density function $\rho(r)$:

$$\rho(r)=\frac{1}{2\pi R^3}\quad\text{for } -R<z<R,\;\; b<R\,,\tag{26}$$

where

$$R=CA^{1/3},\quad C=1.14\text{ F}\,;$$

then

$$\begin{aligned} f_{\gamma A\to vA}(0)=&+\frac{e}{f_\rho}\,\frac{4\pi f_{\rho v}(0)}{\sigma_v-\sigma_\rho}\,\frac{1}{2}R^2(e^{-\sigma_\rho A/2\pi R^2}-e^{-\sigma_v A/2\pi R^2})\\ &+\frac{e}{f_v}\,\frac{4\pi f_{vv}(0)}{\sigma_v}\,\frac{1}{2}R^2(1-e^{-\sigma_v A/2\pi R^2})\,. \end{aligned}$$

The most extreme case is $\sigma_v\gg\sigma_\rho$; then we obtain

$$\begin{aligned} f_{\gamma A\to vA}(0)=&+\frac{e}{f_\rho}\,\frac{2\pi f_{\rho v}(0)}{\sigma_v-\sigma_\rho}\,C^2A^{2/3}e^{-\sigma_\rho A^{1/3}/2\pi C^2}\\ &+\frac{e}{f_v}\,\frac{2\pi f_{vv}(0)}{\sigma_v}\,C^2A^{2/3}\,. \end{aligned}\tag{27}$$

$A^{2/3}e^{-\sigma_\rho A^{1/3}/2\pi C^2}$ increases by factor of 7 while A varies from $A=1$ to $A=240$. Therefore, a study of the A dependence of the cross section in v photo-

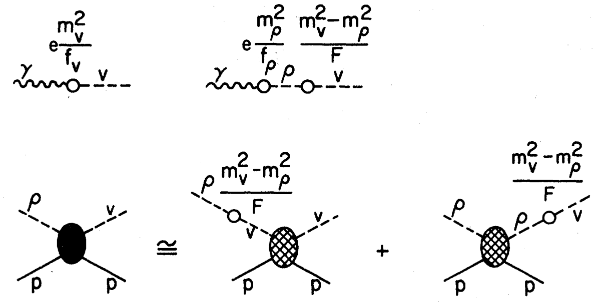

FIG. 5. Diagrams relevant in the weak-coupling model.

production is a very sensitive measurement of the relative strength of the two terms in Eq. (25).

IV. SMALL-COUPLING MODEL

Let us suppose that v is not seen in the experimental search. In such a case we can only place an upper limit on the coupling of γ and ρ with v.

Let us introduce the coupling F and define the coupling between ρ and v to be $(m_v^2 - m_\rho^2)/F$, where we take

$$\frac{m_v^2 - m_\rho^2}{F} \ll \frac{m_\rho^2}{f_\rho}.$$

Then in terms of F we have

$$\frac{f_\rho}{f_v} = \frac{1}{F}\left(1 - \frac{m_\rho^2}{m_v^2}\right),$$

$$\Sigma = \frac{\sigma_v - \sigma_\rho}{F},$$

where we have considered the process only to the first order in $1/F$ as shown in Fig. 5. Furthermore, since σ_v is associated with the area of the proton as seen by v, we expect it to be of the same order of magnitude as that of σ_ρ. The small-coupling model is a special case of the $\Sigma \to 0$ limit. In this model, Eq. (25) becomes

$$f_{\gamma A\to vA}(\theta) = \frac{ik}{F}\frac{e}{f_\rho}\int J_0(kb\theta)\left[1 - 2\exp\left(-\tfrac{1}{2}\sigma_v A\int_{-\infty}^{\infty}\rho(\vec{b}, z')\,dz'\right) + \exp\left(-\tfrac{1}{2}\sigma_\rho A\int_{-\infty}^{\infty}\rho(\vec{b}, z')\,dz'\right)\right] b\,db,$$

where we have also assumed $m_\rho^2 \ll m_v^2$ and $m_v^2 R \ll k$. Note that the A dependence is very sensitive to σ_v/σ_ρ. In fact, if $\sigma_v \approx 10$ mb, destructive interference occurs and there is a dip near $A \approx 65$. Therefore, it is quite important to vary the target nucleus to cover the wide range of A.

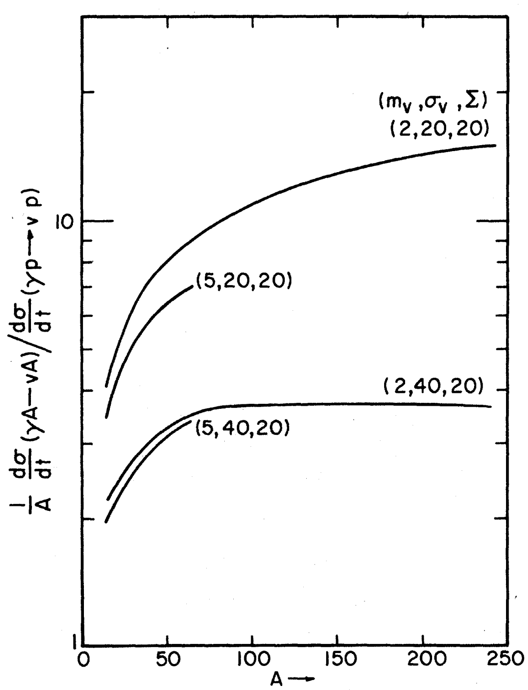

FIG. 6. $(1/A)(d\sigma/dt)(\gamma A \to vA)/(d\sigma/dt)(\gamma p \to vp)$ as a function of A, with $f_\rho/f_v = 0$.

V. RESULTS

We now summarize the predictions of Eq. (15). Since we are only interested in the order-of-magnitude predictions, we will use the cylindrical model with uniform density for the nucleus. This is given in Eq. (26). This model is highly unrealistic, and the results deviate from a more realistic model when $q^{\parallel}R$ becomes sufficiently large. We restrict ourselves to the kinematical region such that the deviation from a more realistic model is small. Also, if there is any deviation, this model will tend to underestimate the v production.

(i) To answer the questions of the absorption effect raised in the Introduction we have plotted

$$\frac{1}{A}\frac{d\sigma}{dt}(\gamma A \to vA)\Big/\frac{d\sigma}{dt}(\gamma p \to vp)$$

as a function of A in Fig. 6. We have set $\Sigma = 20$ mb, $f_\rho/f_v = 0$, and considered various m_v and σ_v. The behavior shown in Fig. 6 also holds for other values of the parameters. Therefore, the absorption effect similar to that of ρ production occurs for v production.

(ii) The question of the interference effect was answered in Sec. III in the small-Σ limit. The fact

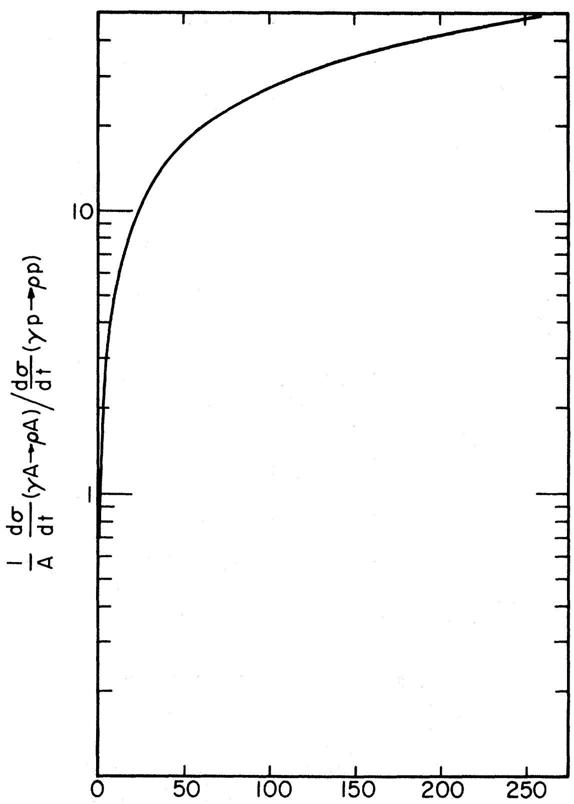

FIG. 7. $(1/A)(d\sigma/dt)(\gamma A\to\rho A)/(d\sigma/dt)(\gamma p\to\rho p)$ as a function of A, at E_γ = 200 GeV.

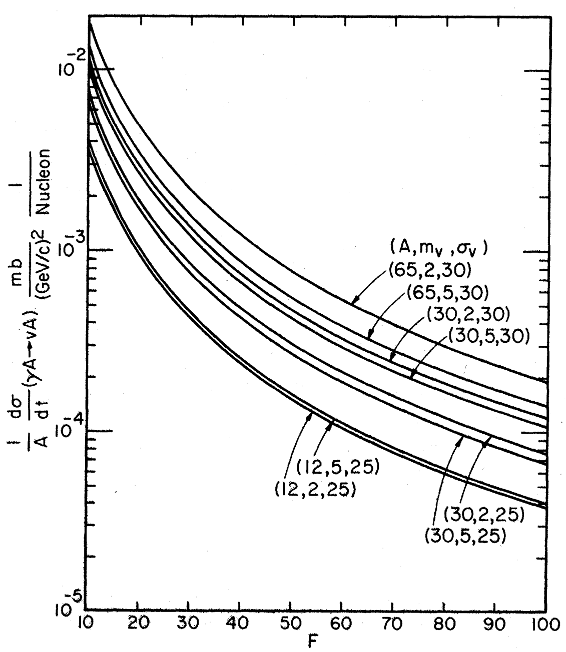

FIG. 8. $(1/A)(d\sigma/dt)(\gamma A\to vA)$ as a function of F; m_v in GeV and σ_v in mb.

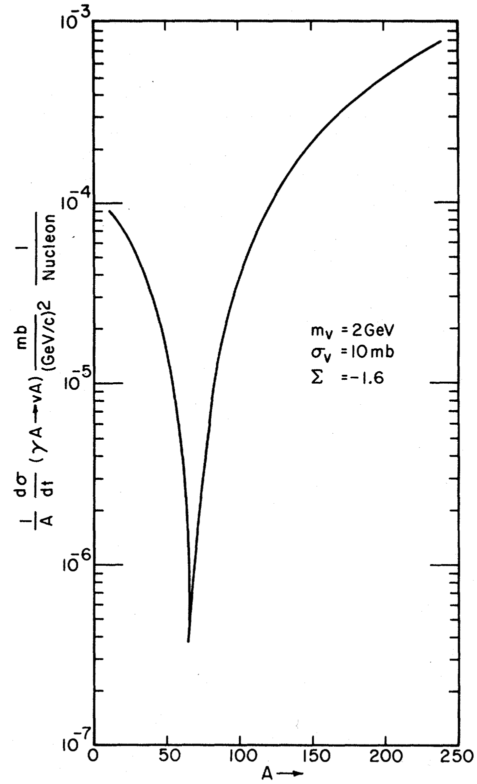

FIG. 9. $(1/A)(d\sigma/dt)(\gamma A\to vA)$ as a function of A; see text.

that v and ρ can be transformed into each other by diffractive scattering will decrease the cross section considerably. To show this we have plotted

$$\frac{1}{A}\frac{d\sigma}{dt}(\gamma A\to\rho A)\Big/\frac{d\sigma}{dt}(\gamma p\to\rho p)$$

as a function of A in Fig. 7. This is to be compared with the behavior shown in Fig. 6. Note that the interference effect also changes the A dependence.

(iii) Let us now be pessimistic and assume that no evidence for the v meson is found. Then the small-coupling model of Sec. IV is applicable. In Fig. 8 we have plotted

$$\frac{1}{A}\frac{d\sigma}{dt}(\gamma A\to vA)$$

as a function of F for various (A, M_v, σ_v). The upper limit on F can be read off considering the experimental limitations. Although we expect $\sigma_v\approx\sigma_\rho$, and therefore the variation of σ_v in Fig. 8 should be sufficient, we have shown

$$\frac{1}{A}\frac{d\sigma}{dt}(\gamma A\to vA)$$

as a function of A for σ_v = 10 mb in Fig. 9 in order to illustrate the effect due to interference between various diagrams shown in Fig. 4. The F depen-

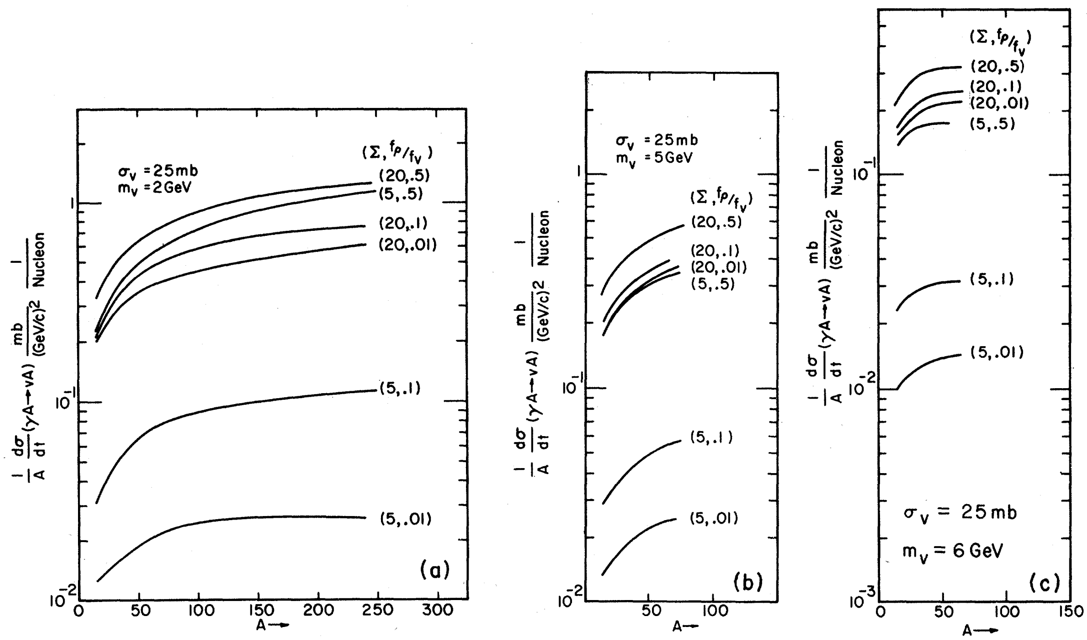

FIG. 10. $(1/A)(d\sigma/dt)(\gamma A \rightarrow vA)$ as a function of A. (a) $m_v = 2$ GeV, (b) $m_v = 5$ GeV, (c) $m_v = 6$ GeV.

dence in Fig. 9 is similar to the F dependences shown in Fig. 8. The origin of the dip is discussed in Sec. IV. It is, therefore, quite important to consider the photoproduction from various nuclei in case of a null result.

(iv) For reference purposes, we have presented the results for various values of the parameters. They are shown in Fig. 10. Examination of these results shows that, if v is found, the A dependence of

$$\frac{d\sigma}{dt}(\gamma A \rightarrow vA)$$

can be used to obtain the parameters f_ρ/f_v and σ_v. On the other hand, the A dependence is quite insensitive to Σ.

Note added. Recently a $JPC = 1--$ particle has been claimed to be observed in $K_L K_S$ resonance in the $p\bar{p}$ reaction [A. Benvenuti *et al.*, Phys. Rev. Letters 27, 283 (1971)]. It will be quite interesting to look for this resonance in the photoproduction.

ACKNOWLEDGMENTS

We would like to thank Professor W. Lee for discussion on the experimental aspects of the problem. One of the authors (A.I.S.) wishes to thank Professor T. D. Lee for getting him interested in the subject and for suggesting the possibility of the small-coupling model.

[1] N. Hicks *et al.*, Phys. Letters 29B, 602 (1969); P. Mostek *et al.*, Phys. Rev. Letters 23, 718 (1969).

[2] It is quite interesting to consider the possibility that there are vector mesons whose conventional quantum numbers $JPC = 1--$ are the same as those for a photon, but they do not couple to the photon, owing to an as yet unknown selection rule. For example, according to the quark model, the recently discovered $R_\alpha(1632)$ has quantum numbers $JPC = 1--$.

[3] If there is no absorption, the rate for the coherent ρ photoproduction on various nuclei should behave as A^2, where A is the atomic number of the nucleus. On the other hand, if there is strong absorption, only those ρ which are created at the back face of the nucleus are seen. For this case, the rate behaves as $A^{4/3}$. Experimentally, the photoproduction of ρ on various nuclei behaves as A^n, where $2 > n > \frac{4}{3}$.

[4] In this paper, we are concerned only with the photoproduction of isovector particles. If the leptonic decay mode of the new vector mesons is detected, for example, e^+e^-, then such an experiment will be sensitive to the isoscalar vector mesons.

[5] This experiment has been approved at NAL: W. Lee, M. Tannenbaum, D. Yount, A. Wattenberg, T. O'Halloran, M. Gormley, and L. Read, NAL report (unpublished).

[6] The approximations used here are discussed by R. Glauber, in *Lectures in Theoretical Physics,* edited by W. E. Brittin and L. G. Dunham (Interscience, New York, 1959), Vol. 1.

[7] After the work was completed, it was pointed out to the authors that the coupled-channel formalism developed in this paper has also been developed by M. Ross and L. Stodolsky, Phys. Rev. Letters 17, 563 (1966); S. J.

Brodsky and J. Pumplin, Phys. Rev. 182, 1794 (1969); G. V. Bochmann and B. Margolis, Nucl. Phys. B14, 609 (1969).

[8]If the v meson exists, the determination of σ_ρ must in principle be modified. But we have seen that ρ production, and therefore σ_ρ, is quite insensitive to the existence of v. We therefore take $\sigma_\rho = 26$ mb.

[9]K. Kölbig and B. Margolis, Nucl. Phys. B6, 85 (1968).