PHYSICAL REVIEW D VOLUME 17, NUMBER 1 1 JANUARY 1978

# Vector-meson-exchange and unitarity effects in low-energy photoproduction*

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(Received 17 August 1977)

Vector-meson-exchange effects are investigated for photoproduction in the Δ region. An improved analysis using a current-algebra model with unitarity corrections and the most recent multipole analysis indicates the necessity of vector-meson contributions. The required vector-meson coupling constants are consistent with quark-model values.

## I. INTRODUCTION

In a previous paper[1] we presented a simple pole model for single-pion photoproduction based on a chiral Lagrangian. The amplitude consisted of the Born terms of the axial-vector ($A$) coupling theory and exchange diagrams for the Δ resonance treated in the zero-width approximation. The presence of the axial-vector Born terms ensured that the low-energy current-algebra results were satisfied. The Δ contributions contained two unknown parameters ($Z$, $Y$) related to the arbitrariness in the choice of Δ couplings. Once these parameters were determined by comparing the predictions of the model with the "experimental results" given by a direct multipole analysis,[2] we were able to account for the essential features of the $l=0$ and $l=1$ multipoles from threshold to about $E_\gamma = 500$ MeV.

While simplicity was one of the most attractive features of the model, it also led to the following shortcomings:

(i) The quantitative agreement with experiment for several multipoles, especially the isoscalar multipole $E_{0^+}(0)$ given by the axial-vector Born terms alone, was not particularly good in the first resonance region except right at threshold.

(ii) The value $Z = 0 \pm \frac{1}{4}$ obtained only barely overlapped the value $Z = -0.45 \pm 0.20$ required by a detailed pole-model treatment of $\pi N$ elastic scattering.[3]

(iii) Finally, because the pole-model amplitudes are real, they cannot reproduce the resonant multipoles $E_{1^+}(\frac{3}{2})$ and $M_{1^+}(\frac{3}{2})$ near the resonance position.

Induced by a recent and significantly improved multipole analysis,[4] we attempt in this paper to remedy the situation by modifying our original model in two ways:

(1) We enlarge the effective Lagrangian in order to include $t$-channel ($\rho$, $\omega$) vector-meson exchanges.

(2) By means of a recently developed technique[5] which uses the experimental elastic (3, 3) phase shift and pole-model results for $\pi N$ elastic scattering[3] we unitarize the resonant multipoles derived from the effective Lagrangian. A similar method is applied to the nonresonant multipoles, although here the effects are less important.

We find that the agreement between the predictions of the modified model,

$$A(\gamma N \to \pi N) = A^{B,A} + A_\Delta(Z, Y) + A_\rho + A_\omega + (\text{unitarity corrections}),$$

and experiment is markedly improved over that of the simpler model. The Δ-coupling parameter set required for the best fit is

$$Z = -0.29 \pm 0.10,$$

$$Y = 0.78 \pm 0.30,$$

which includes a value of $Z$ which is more consistent with the $\pi N$ result. The vector-meson coupling constants are consistent with quark-model values and are close to the vector-meson-dominance predictions.

In Sec. II we describe the enlarged effective Lagrangian and calculate the $l = 0, 1$ multipoles. We carry out the multipole unitarization in Sec. III. In Secs. IV and V we compare the pole-model multipoles with experiment and state our conclusions regarding vector-meson and unitarization effects in low-energy photoproduction.

## II. THE EFFECTIVE LAGRANGIAN AND POLE-MODEL AMPLITUDES

The effective Lagrangian we adopt for our enlarged pole model contains the following interaction terms[6]:

$$\mathcal{L}=[\mathcal{L}_{\pi NN}+\mathcal{L}_{\gamma NN}+\mathcal{L}_{\gamma\pi\pi}+\mathcal{L}_{\gamma NN\pi}+\mathcal{L}_{\Delta\pi N}+\mathcal{L}_{\Delta\gamma N}] +[\mathcal{L}_{\rho NN}+\mathcal{L}_{\rho\gamma\pi}+\mathcal{L}_{\omega NN}+\mathcal{L}_{\omega\gamma\pi}]\,. \tag{1}$$

The part in the first set of brackets is the chiral-invariant Lagrangian on which our previous pole-model treatment was based. The various terms are

$$\mathcal{L}_{\pi NN}=i\left(\frac{f}{\mu}\right)\bar{N}\gamma_\mu\gamma_5\tau^a N\partial_\mu\pi^a\,,$$
$$\mu=\text{charged-pion mass}\,,\quad \frac{f^2}{4\pi}=0.079\pm0.001\,; \tag{2}$$

$$\mathcal{L}_{\gamma NN}+\mathcal{L}_{\gamma\pi\pi}+\mathcal{L}_{\gamma NN\pi} =ie\bar{N}\gamma_\mu\left(\frac{1+\tau^3}{2}\right)NA_\mu+\left(\frac{e}{4M}\right)\bar{N}(\kappa^S+\kappa^V\tau^3)\sigma_{\alpha\beta}F_{\alpha\beta}N +e\epsilon^{ab3}\partial_\mu\pi^a\pi^b A_\mu+i\left(\frac{ef}{\mu}\right)\bar{N}\epsilon^{ab3}\pi^a\tau^b\gamma_\mu\gamma_5 NA_\mu\,, \tag{3}$$

$$M=\text{nucleon mass}=6.722\,\mu\,,\quad \frac{e^2}{4\pi}=\frac{1}{137}$$

$$\kappa^S=\tfrac{1}{2}(\kappa_p+\kappa_n)=-0.060\,,\quad \kappa^V=\tfrac{1}{2}(\kappa_p-\kappa_n)=1.853\,;$$

$$\mathcal{L}_{\Delta\pi N}=g_\Delta\bar{\Delta}^a_\mu\{\delta_{\mu\nu}+[\tfrac{1}{2}(1+4Z)A+Z]\gamma_\mu\gamma_\nu\}N\partial_\nu\pi^a+\text{H.c.}\,, \tag{4}$$

$$\mathcal{L}_{\Delta\gamma N}=-eC\bar{\Delta}^3_\mu\{\delta_{\mu\nu}+[\tfrac{1}{2}(1+4Y)A+Y]\gamma_\mu\gamma_\nu\}\gamma_\lambda\gamma_5 NF_{\nu\lambda} +\text{H.c.}\,.$$

In the above $\Delta$ interactions the arbitrary parameter $A$ is related to the "nonpole" terms in the general $\Delta$ propagator[7] but does not appear in the photoproduction amplitude. Values for the unknown "off-mass-shell coupling parameters" $Z$ and $Y$, and the $\Delta$ coupling constant $C$ are determined by comparing the pole model with experiment.

The additional vector-meson interactions are given by the following terms which result in a breakdown of partial conservation of axial-vector current:

$$\mathcal{L}_{VNN}=ig_{1V}\bar{N}\gamma_\mu NV_\mu+\tfrac{1}{2}g_{2V}\bar{N}\sigma_{\mu\nu}NV_{\mu\nu}\,,$$
$$\mathcal{L}_{V\gamma\pi}=i\left(\frac{\lambda_V}{4\mu}\right)\epsilon_{\mu\nu\lambda\sigma}F_{\mu\nu}V_{\lambda\sigma}\pi\,, \tag{5}$$
$$V=(\rho,\omega)\,,\quad V_{\lambda\sigma}=\partial_\lambda V_\sigma-\partial_\sigma V_\lambda\,,$$

where isospin indices are omitted.

The coupling constant $\lambda_\omega$ can be determined directly from the experimental decay with $\Gamma(\omega\to\pi\gamma)$, with the result[8]

$$\lambda_\omega=(0.357\pm0.011)e\,. \tag{6}$$

For $\lambda_\rho$ we shall take

$$\lambda_\rho=\tfrac{1}{3}\lambda_\omega\,, \tag{7}$$

which follows[9] from the simple nonrelativistic quark model with ideal vector-meson mixing.[10] A recent measurement[11] of $\Gamma(\rho\to\pi\gamma)$ corresponds to a considerably smaller $\lambda_\rho$ than that given by (6) and (7), but, in view of the small amount of data and the subtle nature of this single measurement, we shall rely on the simple quark-model relation (7).

To specify the vector coupling constants $g_{1V}$, we shall again assume maximal symmetry,[10] so that

$$g_{1\omega}=3g_{1\rho}\,, \tag{8}$$

and also $\rho$-meson universality,[12]

$$g_{1\rho}=\tfrac{1}{2}f_\rho\,, \tag{9}$$

where $f_\rho$ is the usual $\gamma$-$\rho$ coupling. Using the value[8] $f_\rho^2/4\pi\simeq2.26$, which is typical of those values extracted from recent Orsay storage-ring data, we obtain

$$g_{1\rho}\simeq2.66\,. \tag{10}$$

We leave the tensor coupling constants $g_{2V}$, or rather the ratios $(g_{2V}/g_{1V})$, as adjustable parameters to be determined by fitting the data. We expect to obtain values consistent with the predictions of simple vector-meson dominance[12]:

$$\frac{g_{2\rho}}{g_{1\rho}}=\frac{\kappa^V}{M}\simeq0.28\,,\qquad \frac{g_{2\omega}}{g_{1\omega}}=\frac{\kappa^S}{M}\simeq-0.009\,. \tag{11}$$

We now proceed to use the effective Lagrangian to calculate to second order in perturbation theory the amplitude for the photoproduction process,

$$\gamma(k)+N(p_1)\to\pi^a(q)+N(p_2)\,.$$

The resulting axial-vector Born terms and the $\Delta$-exchange contributions are given in terms of the usual Chew-Goldberger-Low-Nambu invariant amplitudes by Eqs. (12)–(17) of Ref. 1.

The $\rho$ and $\omega$ exchanges produce the following contributions to the isoscalar and isospin-even amplitudes, respectively:

$$A_1(0)=\frac{-(\lambda_\rho/\mu)g_{2\rho}t}{m_\rho^2-t}\,,\quad A_2(0)=\frac{(\lambda_\rho/\mu)g_{2\rho}}{m_\rho^2-t}\,,$$
$$A_4(0)=\frac{(\lambda_\rho/\mu)g_{1\rho}}{m_\rho^2-t}\,,$$
$$A_1(+)=\frac{-(\lambda_\omega/\mu)g_{2\omega}t}{m_\omega^2-t}\,,\quad A_2(+)=\frac{(\lambda_\omega/\mu)g_{2\omega}}{m_\omega^2-t}\,, \tag{12}$$
$$A_4(+)=\frac{(\lambda_\omega/\mu)g_{1\omega}}{m_\omega^2-t}\,,$$

for which we take $m_\rho^2\simeq m_\omega^2\simeq30\mu^2$.

It should be emphasized that we are not regarding these vector-meson contributions as dispersion-theory pole terms which would prescribe the change $t \to m_V^2$ in the numerators of $A_1(0, +)$, but rather we obtain these contributions by means of a consistent treatment of the entire effective Lagrangian in the tree approximation. As a result, our predictions for the threshold multipoles $E_{0+}(0, +)$ agree well with experiment. Neglecting the small $\Delta$ contributions at threshold and using the vector-meson-dominance ratios (11), the tree approximation (TA) predicts

$$\begin{aligned} 1000E_{0^+}(0)_{\rm TA} &= -1.40\,\mu^{-1}, \\ 1000E_{0^+}(+)_{\rm TA} &= -0.95\,\mu^{-1}, \end{aligned} \tag{13}$$

while typical experimental values[13] are

$$\begin{aligned} 1000E_{0^+}(0)_{\rm exp} &= (-1.40 \pm 0.50)\,\mu^{-1}, \\ 1000E_{0^+}(+)_{\rm exp} &= (-0.60 \pm 0.90)\,\mu^{-1}. \end{aligned} \tag{14}$$

If we use instead the dispersion-theory (DT) pole terms for the nucleon and the vector mesons, we obtain

$$\begin{aligned} 1000E_{0^+}(0)_{\rm DT} &= -2.93\,\mu^{-1}, \\ 1000E_{0^+}(+)_{\rm DT} &= -6.14\,\mu^{-1}, \end{aligned} \tag{15}$$

which disagree with experiment. Radutskii and Serdyutskii[14] have correctly observed that any model for photoproduction at threshold which includes dispersion-theory pole terms for the vector mesons must also include a mechanism which almost cancels them out. In our approach, agreement with experiment is achieved by the combination of much smaller vector-meson contributions and axial-vector nucleon Born terms rather than the pseudoscalar Born terms of dispersion theory.

In summary, our enlarged pole model for photoproduction, expressed in terms of invariant amplitudes is given by the sum of the axial-vector Born terms and $\Delta$ contributions given in Ref. 1, and the vector-meson contributions (12)

$$A_i = A_i^{B,A} + A_i^{\Delta} + A_i^{\rho} + A_i^{\omega}.$$

To compare our model with experimental results given by a direct multipole analyses we now calculate the real parts of the $l = 0$ and $l = 1$ multipoles from the invariant amplitudes using the projection formulas [Eqs. (21)–(24)] listed in Ref. 1. Then the final step is to unitarize the resonant multipoles as described in the next section of this paper.

## III. MULTIPOLE UNITARIZATION

Multipoles of definite $s$-channel isospin, spin, and parity are required by unitarity to have the same complex phase as the corresponding amplitude in the elastic regime.[15] In the $\Delta$ region the elastic phase shifts are all small with the exception of the (3, 3) phase shift $\delta$. Thus, all of the photoproduction multipoles will be nearly real except for $M_{1^+}(\frac{3}{2})$ and $E_{1^+}(\frac{3}{2})$ which resonate. For all of the nearly real multipoles we expect that unitarity effects are small and our model can be almost directly compared with the real parts of the multipoles. For the resonant multipoles we must be more careful. A simple insertion of a width factor into the $\Delta$-resonance denominators of the resonant-invariant amplitudes will generally violate Watson's theorem.[15] The reason for this is that the nonresonant amplitudes, primarily the Born terms, have nonzero projections into the resonant $M_{1^+}(\frac{3}{2})$ or $E_{1^+}(\frac{3}{2})$ multipoles. Thus, even though the $\Delta$ term may be given the correct phase, the phase of the total multipole will not be $\delta$.

The resonant multipoles can be calculated using our effective Lagrangian by means of the experimental (3, 3) phase shift, unitarity, and the following trick[5]: In both the elastic (3, 3) partial-wave and the resonant multipoles $M_{1^+}(\frac{3}{2})$ and $E_{1^+}(\frac{3}{2})$, the lowest-order Lagrangian calculation yields a pole on the energy axis. In a more complete calculation, decay channel effects would shift the pole position. For our present purpose it is sufficient to assume that the resonance factorizes so that unitarity effects in the effective resonance propagator will be the same for elastic scattering and photoproduction. The unknown resonance pole can thus be eliminated by use of the experimental (3, 3) phase shift $\delta$. The resonant multipole $f_{12}$ can then be expressed[5] as

$$qf_{12} = Ne^{i\delta}\sin(\delta + \delta_p - \delta_e), \tag{16}$$

where $N = v_1/v_2$ is the ratio of the $\Delta$ formation and decay vertices for the multipole and $\delta_e$ and $\delta_p$ are the nonresonant background phase shifts from elastic scattering and photoproduction, respectively.

The (3, 3) phase shift $\delta$ is shown in Fig. 1 by a smooth parameterization of the data.[16] The nonresonant background in elastic scattering is obtained by projecting all exchange contributions from a pole-model treatment of $\pi N$ scattering,[3] except the direct-channel $\Delta$ contribution, into the (3, 3) state. The unitarized phase shift $\delta_e$ is shown[17] in Fig. 2. As the figure indicates, the elastic background is dominated by the axial-vector Born term with the remaining $\sigma$-exchange and $u$-channel $\Delta$-exchange contributions amounting to about 25%. Similarly, the photoproduction nonresonant backgrounds calculated from the effective Lagrangian are shown in Fig. 3. In this case the photoproduction phase $\delta_p$ is taken to be[5]

$$qf_{12}^{\rm back} = N\sin\delta_p\,. \tag{17}$$

The photoproduction backgrounds are again primarily due to the axial-vector Born terms with smaller contributions from $u$-channel $\Delta$ exchange

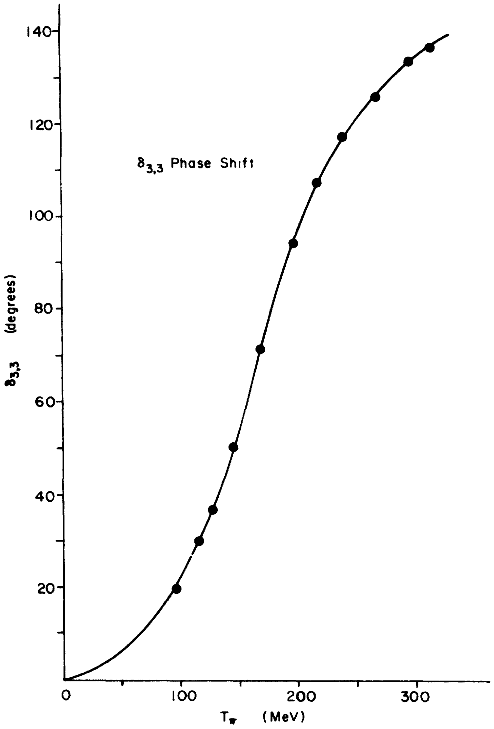


FIG. 1. Experimental (3, 3) $\pi N$ elastic phase shift using the data of Ref. 16.

and $\omega$ exchange.[18]

Once the backgrounds have been calculated, Eq. (16) provides the $M_{1^+}(\frac{3}{2})$ and $E_{1^+}(\frac{3}{2})$ multipole predictions in terms of the $\Delta\gamma N$ coupling constant $C$ which appears multiplicatively in $N$. The determination of $C$ is simplified because the phase of $M_{1^+}(\frac{3}{2})$ and $E_{1^+}(\frac{3}{2})$ is just the (3, 3) elastic phase so that we need only consider the real parts of these multipoles. The ratio of $E_{1^+}$ to $M_{1^+}$ is uniquely specified. Thus, if $C$ is fixed by one point of the $M_{1^+}$ curve, the entire $M_{1^+}$ multipole energy dependence as well as the scale and energy dependence of $E_{1^+}$ is determined. The value of $C$ obtained will be discussed in Sec. IV.

From Figs. 2 and 3 we see that the elastic and $M_{1^+}(\frac{3}{2})$ backgrounds are very similar. As a consequence the $M_{1^+}$ multipole given by the unitarity [Eq. (16)] should be nearly proportional to the elastic (3, 3) amplitude. On the other hand, the $E_{1^+}(\frac{3}{2})$ background is much larger than the elastic background, so by Eq. (16) the $E_{1^+}(\frac{3}{2})$ resonance peak will be shifted significantly to lower energy.

In a similar manner we can unitarize the nonresonant multipoles. If there is no resonance in a given partial wave, $\delta = \delta_e$ and Eqs. (16) and (17) imply that the model amplitude should be multiplied by the elastic amplitude phase factor. The real parts of the nonresonant multipoles are then modified by cos $\delta_e$, which is nearly unity in most cases.

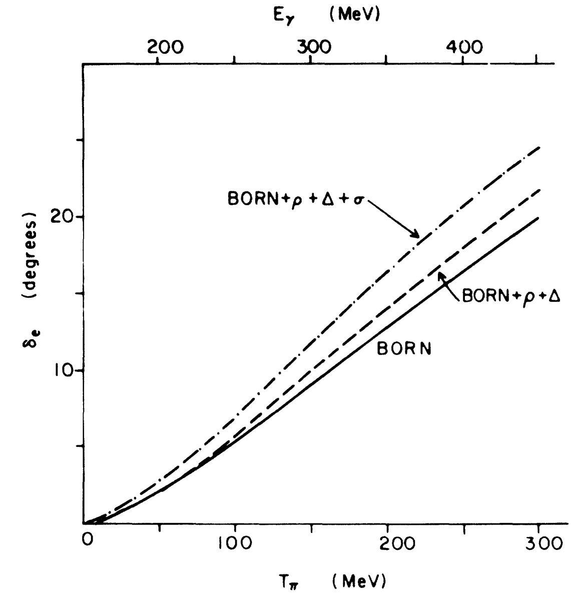


FIG. 2. The elastic background phase.

The largest effect is in the $I = \frac{3}{2}$ $E_{0^+}$ multipole where $\delta_e \simeq -20°$ at the upper energy range, implying a 6% reduction in the real part. Certain process multipoles such as $E_{0^+}(\pi^0 p)$ may have large complex phases as a result of cancellation between the real parts of the component isospin amplitudes. This does not necessarily mean that unitarity effects are large since the phases of the isospin multipoles may be much smaller.

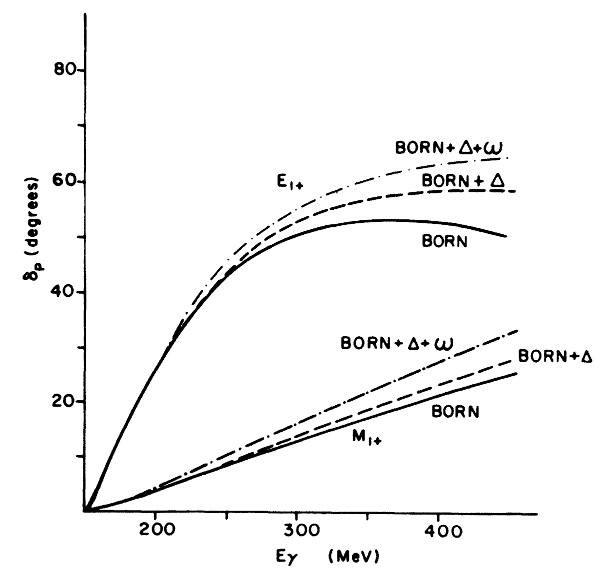


FIG. 3. Photoproduction background phases.

## IV. COMPARISON WITH EXPERIMENT

### A. The data

Several energy-independent multipole analyses[2,4] have been performed which determine the $l=0,1$ multipoles in the $\Delta$ region. The most recent work of Berends and Donnachie[4] incorporated the latest data on $\pi N$ elastic scattering and photoproduction and, as a result, improved significantly on earlier analyses. For our statistical fit we shall rely on the analysis of Berends and Donnachie which covers the energy range $E_\gamma = 250-450$ MeV. However, we shall also display the results of Pfeil and Schwela[2] for energies below 250 MeV and where the two analyses disagree substantially.

### B. The parameters of the model

Our photoproduction model includes the following contributions:

(i) *Axial-vector Born terms* (*no free parameters*).

(ii) $\Delta$ *exchange* (*three parameters*). Recent studies of low-energy $\pi N$ scattering indicate[3,19] that the effective $\Delta$-pole parameters are $M_\Delta = 1220$ MeV, $g_\Delta{}^2/4\pi = 0.27$, and $Z \simeq -0.45$. In photoproduction the $\Delta N\gamma$ coupling introduces additional coupling constants $C$ and $Y$. Since only the combination $Cg_\Delta$ appears in the photoproduction amplitudes,[1] we fix $g_\Delta{}^2/4\pi = 0.27$ and vary the parameters $C$, $Z$, and $Y$. In our previous $\pi N$ work[3] we found that $Z = -0.45 \pm 0.20$, and we anticipate finding a value consistent with this.

(iii) *Vector-meson exchanges* (*two parameters*). The couplings $\lambda_\rho g_{1\rho}$ and $\lambda_\omega g_{1\omega}$ have been fixed at values consistent with experiment, universality, and the quark model. From Eqs. (6), (7), and (10) we find

$$\lambda_\rho g_{1\rho} = 0.096\,, \qquad \lambda_\omega g_{1\omega} = 0.864\,. \tag{18}$$

The ratios $g_{2\rho}/g_{1\rho}$ and $g_{2\omega}/g_{1\omega}$ are treated as free parameters. One might expect these ratios to resemble the vector-meson-dominance values of Eq. (11).

### C. The results of the fit

The five parameters discussed in the previous subsection were varied to optimize agreement between our model amplitudes and the multipoles[4] extracted from experimental data. The resulting parameters values are

$$C = 0.30 \pm 0.01\,, \quad Z = -0.29 \pm 0.10\,, \quad Y = 0.78 \pm 0.30\,, \quad g_{2\rho}/g_{1\rho} = 0.55 \pm 0.20\,, \quad g_{2\omega}/g_{1\omega} = -0.06 \pm 0.10\,. \tag{19}$$

The other coupling constants appearing in our model amplitudes have been held fixed at the values

$$f = 0.996\,, \quad g_\Delta = 1.84\mu^{-1}\,, \quad \lambda_\rho = 0.12e\,, \quad \lambda_\omega = 0.357e\,, \quad g_{1\rho} = 2.66\,, \quad g_{1\omega} = 7.98\,, \tag{20}$$

As has been already discussed, these fixed parameters are either well known from other processes or have strong theoretical foundation.

It might be noted that the fitted values of the $\Delta N\gamma$ coupling constant $C$ coincides exactly with the quark-model expectation.[1] Our present value of $Z = -0.29 \pm 0.10$ is rather sharply determined and agrees better with the $\pi N$ pole-model[3] value $Z = -0.45 \pm 0.20$ than the value $Z = 0 \pm \frac{1}{4}$ required by our simpler photoproduction model[1] with earlier multipole data.[2]

At present, it appears that a state of uncertainty exists as to which values of $Z$ are theoretically preferred. Originally, Nath *et al.*[7] concluded that basic field-theory constraints require the value $Z = \frac{1}{2}$. Recently, their criteria have been reexamined by Jenkins[20] who found that $Z = \pm\frac{1}{2}$ are equally acceptable and, in fact, that, although less preferable, other values of $Z$ are not yet ruled out on theoretical grounds.[21] Although the issue is in doubt from a theoretical viewpoint, our phenomenological approach continues to require a negative value for $Z$.

There is, to our knowledge, no theoretical information regarding the value of $Y$. Our present value of $Y$ is somewhat larger than our previous determination.[1]

The fitted ratio $g_{2\rho}/g_{1\rho} = 0.55 \pm 0.20$ is not quite consistent with the vector-meson-dominance result of 0.28, as given in Eq. (11). Our result agrees well, however, with other analyses, as will be mentioned in Sec. V. The corresponding ratio $g_{2\omega}/g_{1\omega} = -0.06 \pm 0.10$ for $\omega$ exchange is in good accord with the vector-meson-dominance expectation of $g_{2\omega}/g_{1\omega} = -0.009$ given by Eq. (11).

### D. The isospin-$\frac{3}{2}$ multipoles

Displayed in Fig. 4 are the predictions of our model along with the multipole-analyses results of Berends and Donnachie,[4] and Pfeil and Schwela.[2]

Our model agrees quite well with the work of Berends and Donnachie for the nonresonant multipoles $E_{0^+}(\frac{3}{2})$ and $M_{1^-}(\frac{3}{2})$. The more recent results of Berends and Donnachie for $M_{1^-}(\frac{3}{2})$ tend to be substantially more negative than those of Pfeil and Schwela. To a large degree, this change in the data accounts for the preference of $Z \simeq -0.29$ rather than the value $Z \simeq 0$ required to fit the Pfeil and Schwela results in our earlier paper.[1]

Our predictions for the resonant multipoles $M_{1^+}(\frac{3}{2})$ and $E_{1^+}(\frac{3}{2})$, obtained by applying the unitarization procedure described in Sec. III, are independent of the $\rho$ couplings, nearly independent of the off-shell parameters $Z$ and $Y$, and quite insensitive to the choice of the $\omega$ couplings. The model expectations represented by the curves in Fig. 4 show good quantitative agreement for the $M_{1^+}(\frac{3}{2})$ multipole and at least qualitative agreement in the case of $E_{1^+}(\frac{3}{2})$. The turnover of the $E_{1^+}$ curve at and above resonance is not rapid enough to produce the double zero at resonance found by Berends and Donnachie or to provide good agreement at the higher energies. This behavior, however, is quite sensitive to small changes in the background as a 15% background change would increase $\delta_p$ by 30° in the $E_{1^+}$ case.[5]

### E. Photon isovector $I = \frac{1}{2}$ multipoles

Our prediction for the $I = \frac{1}{2}$ isovector multipoles portrayed in Fig. 5 are generally good. For the $E_{0^+}(\frac{1}{2})$ multipole, the data seem to dip below the theoretical prediction in the middle-energy range. In Sec. V we shall see that this is largely accounted for by the isospin (+) multipole.

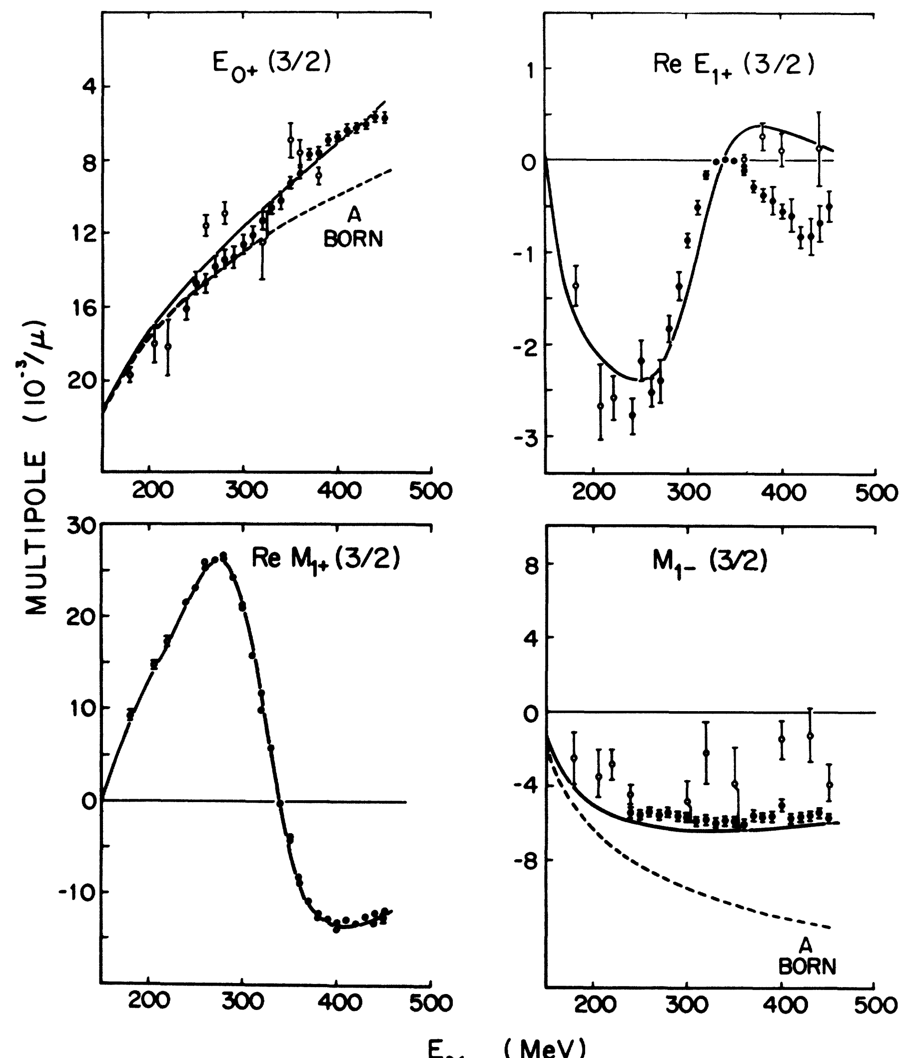


FIG. 4. Isospin-$\frac{3}{2}$ photoproduction mutlipoles. Solid data points are from Ref. 4 and hollow are from Ref. 2. The solid curve represents our model prediction.

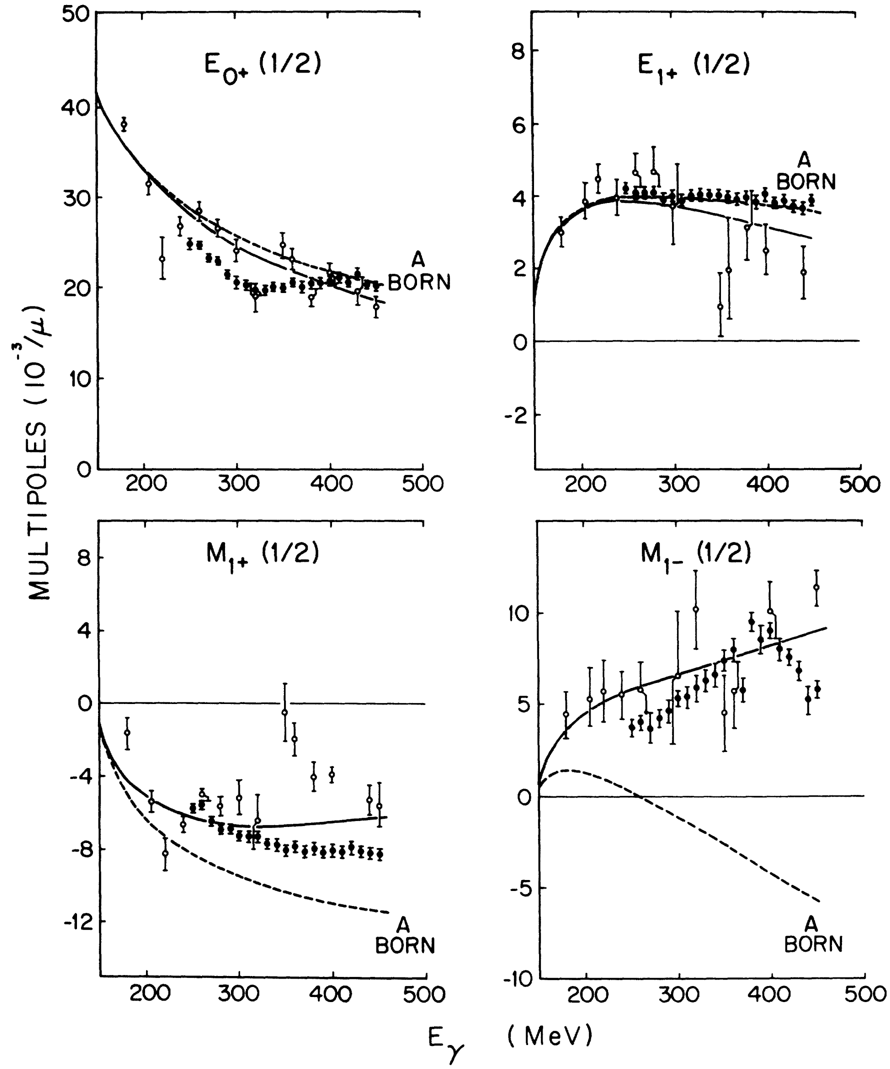


FIG. 5. Isospin-$\frac{1}{2}$ isovector photoproduction multipoles. Data and curves same as Fig. 4.

In the case of $M_{1^+}(\frac{1}{2})$ and $E_{1^+}(\frac{1}{2})$ our predictions tend to deviate slightly from the Berends and Donnachie results toward those of Pfeil and Schwela which are considerably different at many energies. It is interesting to note that while $M_{1^+}(\frac{1}{2})$ and $M_{1^-}(\frac{1}{2})$ require significant $\Delta$ and $\omega$ contributions, the multipole predictions for $E_{0^+}(\frac{1}{2})$ and $E_{1^+}(\frac{1}{2})$ are essentially given by the axial-vector Born terms alone.

### F. The isoscalar multipoles

To complete our description of the low-energy photoproduction data we must consider the isoscalar multipoles. In our model only the axial-vector Born terms and $\rho$ exchange contribute to these multipoles, since $\Delta$ and $\omega$ exchanges are isospin forbidden. As we can see from Fig. 6 the axial-vector Born terms alone are adequate to fit $E_{1^+}(0)$ and $M_{1^+}(0)$, but $E_{0^+}(0)$ and $M_{1^-}(0)$ require significant modification. Fortunately, these latter multipoles receive large contributions from the $\rho$-exchange amplitudes of Eq. (12). As we have mentioned previously the vector-dominance value of $g_{2\rho}/g_{1\rho}=0.28$, given in Eq. (11), lies slightly outside the range of acceptable values given in Eq. (23). The agreement in the case of $E_{0^+}(0)$ is poor for the curve corresponding to $g_{2\rho}/g_{1\rho}=0.30$ which is essentially the vector-meson-dominance prediction.

The importance of the $\rho$ contributions to $E_{0^+}(0)$ and $M_{1^-}(0)$ becomes evident at the higher energies where they are comparable to the axial-vector Born contributions.

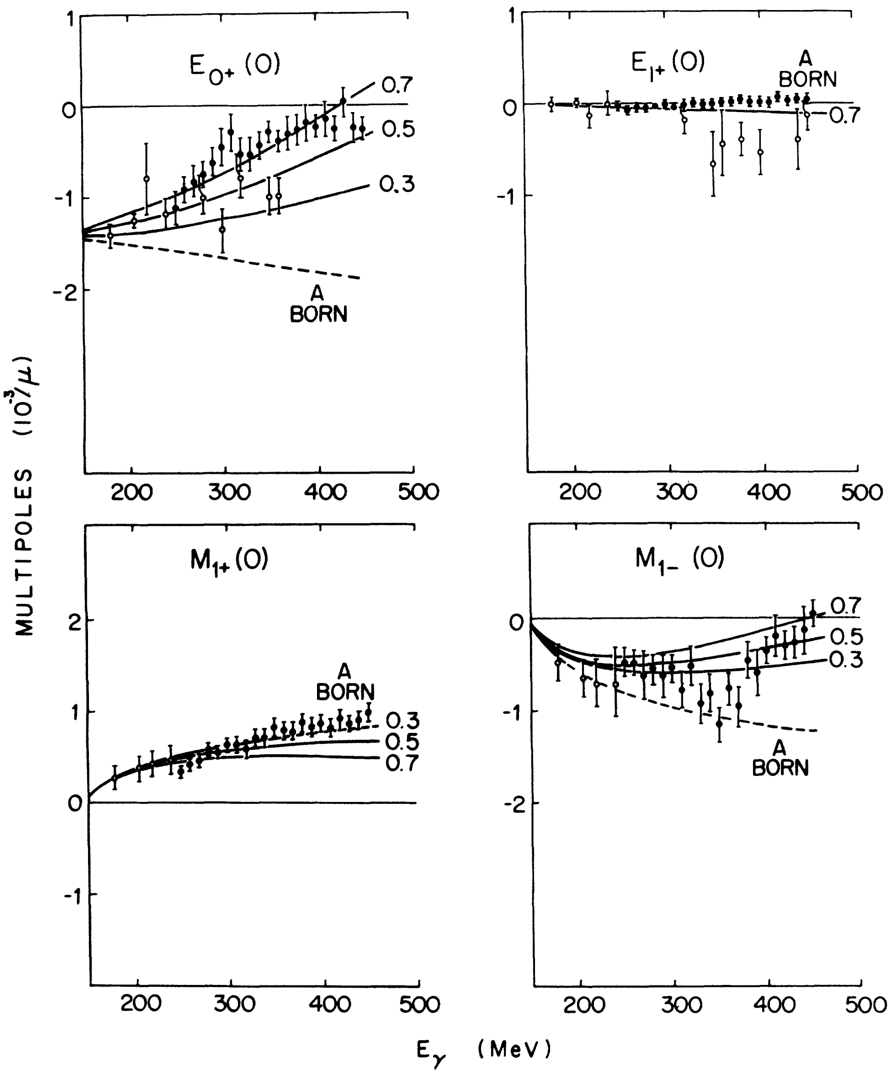


FIG. 6. Isoscalar photoproduction multipoles. Data same as Fig. 4. The curves correspond to different values of $g_{2\rho}/g_{1\rho}$.

## V. DISCUSSION AND CONCLUSIONS

The main theoretical modification of our previous model[1] is the introduction of vector-meson exchanges. The need for $\rho$ exchange in the isoscalar multipoles was discussed in the last section. The corresponding phenomenological need for $\omega$ exchange has not yet been made clear. Since $\omega$ exchange contributes to the (+) amplitudes alone, it is useful to consider the isospin even and odd combinations

$$A(+)=\tfrac{1}{3}[A(\tfrac{1}{2})+2A(\tfrac{3}{2})]\,,\qquad A(-)=\tfrac{1}{3}[A(\tfrac{1}{2})-A(\tfrac{3}{2})]\,. \tag{21}$$

In Fig. 7 we plot the nonresonant multipole combinations $E_{0^+}(\pm)$ and $M_{1^-}(\pm)$. Only axial-vector Born and $\Delta$ exchanges contribute to the $E_{0^+}(-)$ and $M_{1^-}(-)$ multipoles and we see that our model curves fit the data reasonably well.

The effect of varying the $\omega$-coupling ratio $g_{2\omega}/g_{1\omega}$ can be seen in the $E_{0^+}(+)$ and $M_{1^-}(+)$ multipoles of Fig. 7. The $\omega$-exchange contribution is substantial, especially in the case of $M_{1^-}(+)$ where it amounts to 30% of the axial-vector Born contribution. We should also point out that while the representative curves for $g_{2\omega}/g_{1\omega}$ pass through the average of the $E_{0^+}(+)$ data they do not exhibit the structure in the data or the zero at $E_\gamma \simeq 340$ MeV. It is interesting to note that this zero is near the $\Delta$-resonance energy and thus may be spurious.

Some additional discussion of the vector-meson coupling constants is in order. Because of increasingly accurate measurements of the width $\Gamma(\omega\to\pi\gamma)$, the value of $\lambda_\omega$ is becoming well deter-

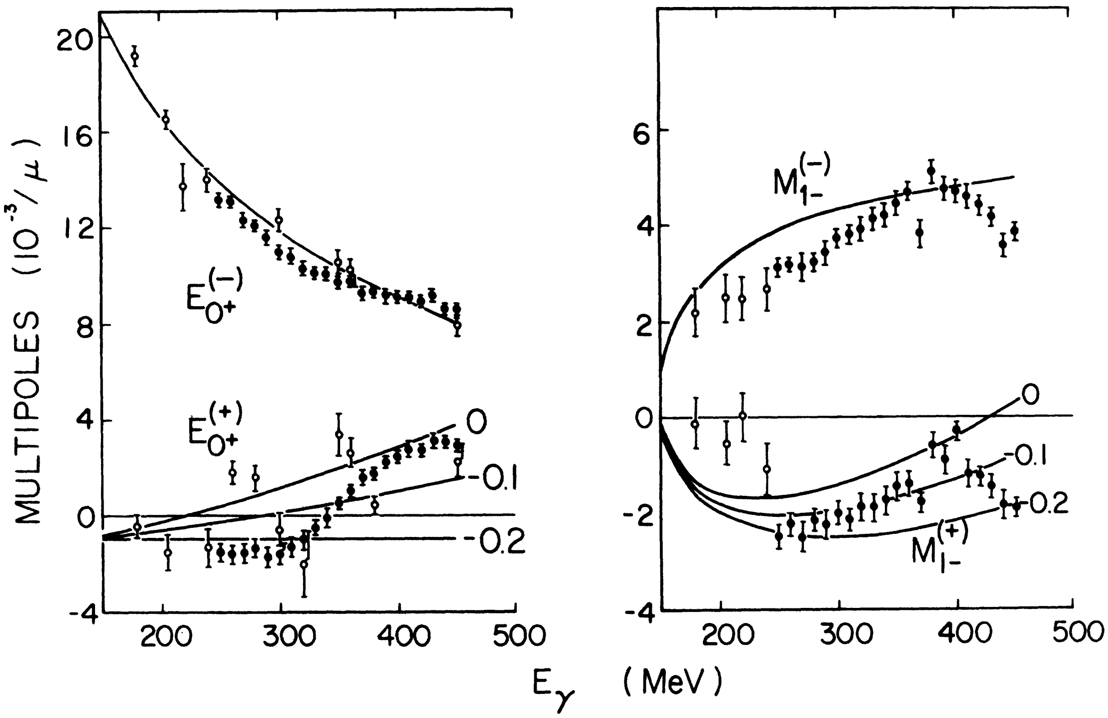


FIG. 7. Isospin-even and -odd photoproduction multipoles. The curves in the even case represent various values of $g_{2\omega}/g_{1\omega}$.

mined. On the other hand, the single existing measurement[11] of width $\Gamma(\rho\to\pi\gamma)$ corresponds to a value of $\lambda_\rho$ approximately one half those predicted by various symmetry schemes[22] for the radiative decays of vector mesons. There is clearly a need for additional measurements to confirm whether or not our decision to use the quark-model relation (7) is justified.

The situation with the *VNN* coupling constants $g_{1V}$ and $g_{2V}$ is also somewhat confused. Even if the universality assumption[9] is valid, the various quoted values of $f_\rho$ depend upon the method used to extract it from $\rho\to e^+e^-$ data and upon the type of finite width corrections applied. A detailed study of the $\rho NN$ vertices in vector-dominance models by Höhler and Pietarinen[23] predicts the value[24]

$$g_{1\rho}=2.63\,, \tag{22}$$

which agrees well with our assumed value $g_{1\rho}\sim 2.66$. They also obtained the result $g_{2\rho}/g_{1\rho}=0.49$; again consistent with our value $g_{2\rho}/g_{1\rho}=0.55\pm0.20$. In a related analysis of the electromagnetic nucleon form factors, Genz and Höhler[25] calculated values of $g_{1\omega}$ and the corresponding vector coupling $g_{1\phi}$ for the $\phi$ meson which, when combined with the Höhler and Pietarinen value (22) of $g_{1\rho}$, are compatible with SU(3), although strong violation of the Zweig-Iizuka rule is implied. As one would expect, their value[25] of

$$g_{1\omega}=17.4\pm2 \tag{23}$$

is quite different than the value $g_{1\omega}\simeq 8$ we calculated, assuming validity of the Zweig-Iizuka rule. It should be emphasized that within the context of our model it would be impossible to fit the isospin-even multipoles using the value (23) without invoking a totally unrealistic ratio $g_{2\omega}/g_{1\omega}$.

The present pole-model analysis of photoproduction leads us to a number of conclusions:

(1) As in the elastic $\pi N$ case,[3] the phenomenological application of $\Delta$-exchange terms favors negative values of the off-shell $\Delta\pi N$ coupling parameter; specifically,

$$Z=-0.29\pm0.10\,.$$

(2) The available low-energy photoproduction data as summarized by direct multipole analyses cannot be accounted for completely by axial-vector Born terms and $\Delta$-exchange contributions alone.

(3) The addition of $\rho$- and $\omega$-exchange contributions, given in terms of coupling constants consistent with the measured width $\Gamma(\omega\to\pi\gamma)$, $\rho$-meson universality, and the quark model, results in a vastly improved pole model.

(4) If supplemented by elastic (3, 3) phase shift data and $\pi N$ pole model results, the resonant

multipoles calculated from the enlarged pole model can be unitarized in a satisfactory manner. The nonresonant multipoles can be similarly unitarized, although in this case the effect is slight.

It should be remarked that coupled with the inclusion of $t$-channel vector-meson exchanges is the risk of double counting. However, we take the point of view that, given our effective Lagrangian approach, the most natural and practical way to calculate the increments to the axial-vector Born and $\Delta$ contributions, required by conclusion (2) above, is to include the exchanges of the lowest-lying vector mesons. An additional motivation is provided by a "hard pion" Lagrangian treatment of photoproduction and electroproduction by Radutsky and Serdyutsky[26] in which vector mesons enter in a natural fashion. It might be noted that vector-meson exchanges are required in electroproduction to account for the nucleon form factors.

We would like to make a few remarks about the uniqueness of our vector-meson-exchange results. It might be argued that nucleon-resonance effects might be incorrectly interpreted as vector-meson exchanges. This seems unlikely because of the resonable structure and coupling constants obtained in our fit; however, let us discuss each case in detail. The $\rho$ exchange only appears in the photon isoscalar amplitude and since the $\Delta(3,3)$ does not couple to this amplitude the only trouble might come from the tail of the $P_{11}(1470)$ resonance which would be manifested in the $M_{1-}(0)$ multipole. Even at its central peak near $E_\gamma = 700$ MeV the $P_{11}(1470)$ is rather difficult to observe in photoproduction. At the upper end of our energy range, $E_\gamma = 450$ MeV, the $P_{11}$ phase shift is only 15° (of which about half is due to nonresonant background[3]), so we do not expect $P_{11}(1470)$ effects to be large. Further suppression in the isoscalar amplitude occurs since the proton and neutron amplitudes cancel to a considerable extent in this case.[27] The $\omega$-exchange effects only occur in the isospin (+) photon isovector amplitudes. Thus by examining the isovector photon isospin (−) amplitude we can fix the off-shell $\Delta(3,3)$ parameters $Y$ and $Z$ independently of the $\Delta$-exchange effects. The result of such a fit gives virtually the same $\Delta$ couplings as in our overall fit, giving additional confidence that the $\omega$ and $\Delta$ contributions have been correctly interpreted.

In summary, properly unitarized pole models for photoproduction and $\pi N$ scattering which incorporate the current-algebra low-energy theorems, presently provide a consistent and reasonably accurate phenomenology for $\pi N$ interactions through the first resonance region. Of course, the validity of this conclusion will be sustained only if these models continue to agree with phase-shift and multipole analyses as these analyses improve in the future.

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*Supported in part by the University of Wisconsin Research Committee with funds granted by the Wisconsin Alumni Research Foundation, and in part by the Energy Research and Development Administration under Contract No. E(11-1)-881, COO-567.


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