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# Photodisintegration of the Deuteron*

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The Chew meson theory has been applied in an attempt to describe the photodisintegration of the deuteron in the region of 100 to 400 Mev, and in particular to explain the appearance of a sort of resonance at about 250 Mev. The process is viewed as a meson photoproduction, followed by meson scattering in deuterium (treated in the impulse approximation) and absorption. The qualitative features of total cross section are roughly reproduced, but the angular distribution does not seem to be fully understood on this model.

FOR the past several years a considerable amount of experimental information has been accumulating concerning the photodisintegration of the deuteron at energies running up to about 500 Mev.[1] These results indicate a resonance of sorts in the total cross section at a photon energy of about 225 Mev. Theoretical calculations have been carried up to about 150 Mev ignoring any explicit meson effects,[2] and somewhat further by means of including a meson magnetic moment.[3] These calculations do not seem to predict any resonant behavior. In view of the observed resonances in both the meson nucleon scattering and meson photoproduction from hydrogen in the $J=\frac{3}{2}$, $T=\frac{3}{2}$ states, one is tempted to suggest that the resonance in the deuteron photodisintegration is due to the scattering in a virtual state of a meson which has been produced by the photon, the meson being finally absorbed by one of the two outgoing nucleons. In order to attempt a quantitative description of the effects of such processes, it is necessary to fix on a particular meson theory to describe the meson-nucleon interaction. Since relativistic forms of meson theory have not been very successful, and since there is at present no reasonably valid approximation method which can be applied to them, it seems advisable to use Chew's form of meson theory.[4] This theory has several advantages. First, it is not a complete theory, and is rather more a phenomenological approach in the sense that no attempt is made to describe a large group of meson phenomena, such as $S$-wave interactions, relativistic effects, or heavy mesons. Of these the $S$-wave interactions are just ignored, and the rest are assumed to be describable by a cutoff on the momentum of any virtual meson, thus restricting one to low energies and providing an extra parameter. The theory is constructed to agree with the important qualitative features observed in the meson-nucleon interaction; namely, strong $P$-wave couplings, pseudoscalar mesons and conservation of isotopic spin. Second, the coupling constant is small, so that there exist fairly reasonable approximation methods. Third, the theory is relatively easy to use, at least compared with the existing relativistic theories. Finally, the theory agrees fairly well with all low-energy $P$-wave meson effects.

For the reasons indicated above, the Chew theory will be used in an attempt to describe the deuteron photodisintegration up to energies of several hundred Mev. Since even a relatively simple theory such as this becomes quite complicated when applied to two-nucleon problems, some rather drastic approximations will of course be necessary. It is therefore not to be expected that more than merely qualitative features of the cross section will be reproduced.

The contribution to the photodisintegration process by explicit meson effects may essentially be viewed as the photoproduction of a meson from one or the other of the nucleons, with the subsequent scattering and absorption of the meson. In the spirit of the impulse approximation,[5] it may be expected that processes which differ from others only in that they involve a large number of meson exchanges between the nucleons will be weaker and can be neglected. That is, it will be assumed that the amplitude for the exchange of mesons is small.

The photoproduction process may, in the Chew theory, take place in essentially three ways. First, an $S$-wave meson may be produced by the three-field coupling $\boldsymbol{\sigma}\cdot\hat{e}$. Second, mesons of any angular momentum may be produced through the meson current interaction. Finally, $P$-wave mesons may be produced by allowing the photon to be absorbed by the nucleon through its anomalous magnetic moment. Once a $P$-wave meson is produced, it may rescatter. No other meson may interact further with the nucleon. At low energies, therefore, there should appear $S$-wave mesons, and $P$-wave mesons, whose amplitude is enhanced by

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[1] J. Keck and A. Tollestrup (private communication); see also J. Keck and A. Tollestrup, preceding paper [Phys. Rev. **101**, 360 (1956); L. Allen Jr., Phys. Rev. **98**, 705 (1955); E. A. Whalin, Phys. Rev. **95**, 1362 (1954); Schriever, Whalin, and Hansen, Phys. Rev. **94**, 763 (1954); Keck, Littauer, O'Neill, Perry, and Woodward, Phys. Rev. **93**, 827 (1954).

[2] J. F. Marshall and E. Guth, Phys. Rev. **78**, 738 (1950); L. I. Schiff, Phys. Rev. **78**, 733 (1950).

[3] Y. Nagahara and J. Fujimura, Progr. Theoret. Phys. (Japan) **8**, 49 (1952).

[4] G. F. Chew, Phys. Rev. **94**, 1748 (1954).

[5] G. F. Chew, Phys. Rev. **80**, 196 (1950); G. F. Chew and G. C. Wick, Phys. Rev. **85**, 636 (1952); G. F. Chew and M. L. Goldberger, Phys. Rev. **87**, 778 (1952).

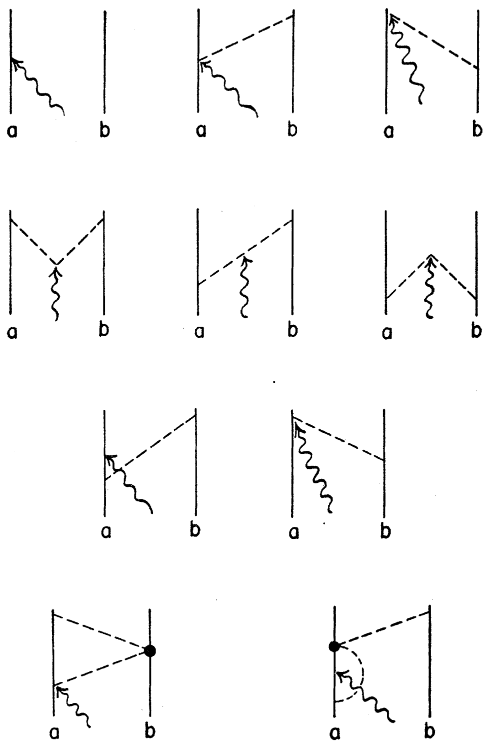


FIG. 1. Diagrams included in the $R$-matrix as given in Eq. (10). Heavy lines represent nucleons, dashed lines mesons, and wavy lines photons. ● represents the scattering of a meson by a nucleon.

rescattering from the nucleon in the resonant $J=\frac{3}{2}$, $T=\frac{3}{2}$ state.

Part of the photodisintegration cross section should therefore be simply the photoproduction of the type described above, from one nucleon, with the photoproduced meson being absorbed by the other nucleon [in order that it may scatter in the $(\frac{3}{2},\frac{3}{2})$ state from the first nucleon]. As stated above, processes in which the meson rescatters from the second nucleon and then returns to the first to be absorbed involve two meson exchanges and will be assumed small compared to the one-meson exchange. However, the presence of the second nucleon allows a further process, which there is no reason for neglecting. The $S$-wave meson produced from the first nucleon may scatter in the $(\frac{3}{2},\frac{3}{2})$ state from the second and be absorbed by the first. This process admittedly involves a two-meson exchange, but here there is no corresponding one-meson exchange process. The only way in which the $S$-wave meson can resonate is by means of a two-meson exchange. In other words, it is not valid to assume that a two-meson exchange with one type of photon coupling is necessarily smaller than a one-meson exchange with another photon coupling; all that is assumed is that a two-meson exchange with one photon coupling is smaller than a one-meson exchange with the same coupling.

With the above as an introduction, we may now proceed to set up the process formally. The interaction Hamiltonian is $V_1+V_2$, where $V_1$ describes the meson-nucleon interactions, and is given by

$$V_1=H_{nm}(a)+H_{nm}(b)$$

for nucleons $a$ and $b$, and $V_2$ describes the electromagnetic interactions, so that $V_2=H_n+H_m+H_3$, denoting the nucleon and meson currents and the three-field interaction, respectively. The $R$ matrix,[6] to lowest order in $V_2$, is then

$$R=V_2+R_1+V_2\frac{1}{a}R_1+R_1\frac{1}{a}V_2+R_1\frac{1}{a}V_2\frac{1}{a}R_1, \tag{1}$$

where

$$a=E_0-H_0+i\epsilon,\quad R_1=V_1+V_1\frac{1}{a}R_1. \tag{2}$$

The desired matrix element of $R$ is one connecting states of 0 mesons and 1 photon to 0 mesons. Since $V_1$ can connect states of $n$ mesons only with states of $n\pm 0$, 1, or 2 mesons, the desired matrix element between two states of 0 mesons is

$$\begin{aligned}(0|R|0)=(0|V_2|0)&+\sum_{0,1,2}(0|R_1|0,1,2)\frac{1}{a}(0,1,2|V_2|0)\\&+\sum_{0,1,2}(0|V_2|0,1,2)\frac{1}{a}(0,1,2|R_1|0)\\&+\sum_n(0|R_1|n)\frac{1}{a}(n|V_2|n\pm 0,1,2)\\&\times\frac{1}{a}(n\pm 0,1,2|R_1|0).\end{aligned} \tag{3}$$

The matrix element $(n|R_1|0)$ includes all Feynman diagrams leading from a state of 0 mesons to one of $n$ mesons under the meson-nucleon interaction. It may be written

$$(n|R_1|0)=(n|R_1|0)'\left[1+\frac{1}{a}(0|R_1|0)\right], \tag{4}$$

where $(n|R_1|0)'$ contains all terms which at no intermediate stage have a state of 0 mesons, and $(0|R_1|0)$ contains, as indicated, all terms going from 0 mesons to 0 mesons. Now, if $V$ denotes the two-nucleon potential derived from this meson theory,

$$(0|R_1|0)=V+V\frac{1}{a}V+V\frac{1}{a}V+\cdots. \tag{5}$$

[6] M. L. Goldberger and M. Gell-Mann, Phys. Rev. **91**, 398 (1953).

Therefore

$$
\begin{aligned}
|\psi) &\equiv |0)+\frac{1}{a}(0|R_1|0)|0) \\
&= |0)+\frac{1}{a}\left(V+V\frac{1}{a}V+\cdots\right)|0) \\
&= |0)+\frac{1}{a}V|\psi), \qquad (6)
\end{aligned}
$$

where $|\psi)$ represents the wave function for two nucleons under the interaction of the potential $V$.[7] It is then possible to write the desired matrix element as

$$
(0|R|0)=\left(\psi_f|\,V_2+R_1'\frac{1}{a}V_2 + V_2\frac{1}{a}R_1'+R_1'\frac{1}{a}V_2\frac{1}{a}R_1'|\psi_i\right), \qquad (8)
$$

where $R_1'$ represents all the meson-nucleon interaction diagrams with no intermediate state of 0 mesons, and $\psi_i$ and $\psi_f$ are the exact (under the meson-nucleon interaction) initial and final states of the two-nucleon system. Thus $\psi_i$ represents the deuteron wave function, and $\psi_f$ the outgoing neutron and proton. From now on it will merely be assumed that meson theory could predict accurate wave functions $\psi_f$ and $\psi_i$, and we shall use phenomenological descriptions of them.

As yet no approximations have been made. At this point we shall assume Chew's form of meson theory,[4] and use the approximations associated with it. Consider, for example, the term $(\psi_f|V_2(1/a)R_1'|\psi_i)$ in Eq. (8). This will have the form of a number of meson-nucleon interactions, separated by energy denominators of the form $(E_0-E_{\text{intermediate}}+i\epsilon)^{-1}$, and followed by a photon interaction. Now $E_0$ is the photon energy, and $E_{\text{intermediate}}$ is, since the photon has not yet been absorbed, $E_0$ plus the energy of whatever mesons are present. The energy denominators can therefore never vanish. On the other hand, in a term like

$$
\left(\psi_f|R_1'\frac{1}{a}V_2|\psi_i\right),
$$

the meson interactions all occur after the photon absorption, so the energy denominators have the form $(E_0-E_{\text{mesons}}+i\epsilon)^{-1}$, and this term may have a pole. Since the coupling constant in the Chew theory is small, one may expect that higher-order processes will be unimportant unless there is some offsetting effect produced by their introduction, such as vanishing energy denominators. Therefore, all terms describing interactions taking place before the photon interaction will be included to the lowest order only.

The matrix element then takes the following form:

$$
\begin{aligned}
(0|R|0)=\Bigg(\psi_f|\,V_2&+H_{nm}^{-}\frac{1}{a}V_2+V_2\frac{1}{a}H_{nm}^{+} \\
&+H_{nm}^{-}\frac{1}{a}V_2\frac{1}{a}H_{nm}^{+}+H_{nm}^{-}\frac{1}{a}R_s\frac{1}{a}\bigg(H_3+H_n\frac{1}{a}H_{nm}^{+} \\
&+H_{nm}^{-}\frac{1}{a}H_m+H_m\frac{1}{a}H_{nm}^{+}\bigg)|\psi_i\Bigg). \qquad (9)
\end{aligned}
$$

The first four terms here represent the lowest order contribution; the first of order $e$, the other three of order $ef^2$. The last part may be interpreted as the photoproduction of a meson [through $H_3+H_n(1/a)H_{nm}^{+}+H_{nm}^{-}(1/a)H_m+H_m(1/a)H_{nm}^{+}$] from either nucleon, which is then scattered ($R_s$ representing the scattering of a meson by a deuteron) and then reabsorbed (through $H_{nm}^{-}$) on either nucleon. Now in the calculation of the direct photoproduction of mesons from protons, it is found that the terms in $H_m$ are much smaller than those in $H_n$.[8] The former will therefore be neglected. To describe the scattering of a meson by a deuteron, that is, to calculate $R_s$, we shall use the impulse approximation. There has been a considerable amount of argument in favor of the validity of this approximation as applied to deuterons[9]; nevertheless, there is no clear justification for its use, and it may tend to overestimate the amplitude significantly.[10] We thus take $R_s=R_s^{(a)}+R_s^{(b)}$, where $R_s^{(a)}$ represents the scattering amplitude in the $(\frac{3}{2},\frac{3}{2})$ state for mesons on nucleon $a$ alone.

The resulting matrix element, in its final form, is

$$
\begin{aligned}
(0|R|0)&=(\psi_f|M|\psi_i) \\
&=\Bigg(\psi_f|\,H_n+H_{nm}^{-}\frac{1}{a}H_3^{+}+H_3^{-}\frac{1}{a}H_{nm}^{+} \\
&\quad+H_{nm}^{-}\frac{1}{a}H_{nm}^{-}\frac{1}{a}H_m+H_{nm}^{-}\frac{1}{a}H_m\frac{1}{a}H_{nm}^{+} \\
&\quad+H_m\frac{1}{a}H_{nm}^{+}\frac{1}{a}H_{nm}^{+}+H_{nm}^{-}\frac{1}{a}H_n\frac{1}{a}H_{nm}^{+} \\
&\quad+H_{nm}^{-}\frac{1}{a}R_s\frac{1}{a}H_3^{+} \\
&\quad+H_{nm}^{-}\frac{1}{a}R_s\frac{1}{a}H_n\frac{1}{a}H_{nm}^{+}|\psi_i\Bigg). \qquad (10)
\end{aligned}
$$

The first bracket here contains the lower-order terms (of order $e$ or $ef^2$); the second contains the "resonant" terms. $R_s$ represents the sum of the scattering amplitude of a meson from the two nucleons. Thus the notation is $R_s=R_s(a)+R_s(b)$, $H_3=H_3(a)+H_3(b)$, etc. The diagrams included here are shown in Fig. 1, dropping

B. Lippmann and J. Schwinger, Phys. Rev. **79**, 469 (1950).

[8] G. F. Chew and F. E. Low (private communication).

[9] G. F. Chew and H. W. Lewis, Phys. Rev. **84**, 779 (1951); Isaacs, Sachs, and Steinberger, Phys. Rev. **85**, 803 (1952); Fernbach, Green, and Watson, Phys. Rev. **82**, 980 (1951).

[10] K. A. Brueckner, Phys. Rev. **89**, 834 (1953).

vertex renormalizations and diagrams contributing to the nucleon anomalous moment. The matrix elements of operators such as $H_{nm}$ are most easily calculated between plane wave states. It will therefore be most convenient to evaluate them for plane wave states, and then integrate over the momentum distribution in $\psi_i$ and $\psi_f$. Explicitly,

$$(\psi_f|M|\psi_i)=\int\frac{d^3\mathbf{P}_f}{(2\pi)^3}\int\frac{d^3\mathbf{P}_i}{(2\pi)^3}\times(\psi_f|\mathbf{P}_f)(\mathbf{P}_f|M|\mathbf{P}_i)(\mathbf{P}_i|\psi_i). \quad (11)$$

A crude approximation suggests itself here. Since the momentum distribution in the deuteron has a half-width of only about 50 Mev, one might expect the important part of the integral over $d^3\mathbf{P}_i$ to come from the region $0<P_i\leq 50$ Mev. This momentum is fairly small compared with the others occurring: (e.g., the momentum of the outgoing particles ranges from 300 Mev for a 100-Mev photon to 650 Mev for a 400-Mev photon). Therefore $\mathbf{P}_i$ could be neglected in $(\mathbf{P}_f|M|\mathbf{P}_i)$, giving the result

$$(\psi_f|M|\psi_i)=\int\frac{d^3\mathbf{P}_f}{(2\pi)^3}(\psi_f|\mathbf{P}_f)(\mathbf{P}_f|M|0)\psi_i(r=0).$$

This expression is really good, of course, only for very high photon energies, and shows the sensitivity of the process to the deuteron wave function near the origin at such energies. For the energy range of interest here, it is unlikely that $P_i$ is negligible, so it will be necessary to investigate more carefully the dependence of the matrix element on the choice of the deuteron wave function.

For the final state wave function a plane wave will be chosen. This is exact for states of odd orbital angular momentum for a 50% exchange force. Since no multipole expansion will be made here, however, this is not necessarily an ideal choice. The initial state must be represented by a phenomenological deuteron wave function. Unfortunately, there are a variety of wave functions, all of which fit the low-energy experimental data fairly well, and all of which predict somewhat different results at high energies. At low energies, only the outer part of the wave function is very important, so that the behavior near the origin is undetermined, but at very high energies this is the important region.

What has been done here is to consider certain of the terms in Eq. (10) [namely, $H_{nm}^-(1/a)H_3+H_3(1/a)H_{nm}^+$] for which the integral over the momentum distribution in the deuteron may be explicitly carried out, and evaluate these for various choices of $\psi_i$. The results indicate a variation in the matrix element of magnitudes up to 30% between square well, Hulthén, and repulsive-core wave functions, the Hulthén giving the largest and the repulsive core the smallest results. It is also found that the approximation of Eq. (12) is excellent for a square-well wave function, but gives too large an answer for a Hulthén well, and too small an answer (naturally) for a repulsive core. This result can easily be understood, since the square-well function is almost flat over most of the volume of importance here.

The procedure will therefore be to assume a square-well wave function and use the approximation of Eq. (12) to insert it into the plane wave matrix elements. The final results may then be expected to be somewhat smaller for a repulsive-core wave function; how much smaller, depends on the size of the core.

Consider first the lowest order effects, as described by the first bracket in Eq. (10). These contain no meson-nucleon scattering processes, so they should not be expected to produce any resonant effects. Through the term in $H_n$, the direct photodisintegration (that is, without meson effects) is included, so that the low-energy behavior will be the same as in previous calculations.

The explicit form of the couplings is[11]

$$H_n=-\sum_{a,b}\frac{ie}{2M_a}\frac{1+\tau_a{}^3}{2}\mathbf{A}\cdot\nabla_a-\tfrac{1}{2}[(\mu_p+\mu_n)+(\mu_p-\mu_n)\tau_a{}^3]\boldsymbol{\sigma}_a\cdot(\nabla\times\mathbf{A}),$$

$$H_{nm}^-=i(4\pi)^{\frac{1}{2}}\frac{f}{\mu}\frac{1}{(2\omega)^{\frac{1}{2}}}[(\boldsymbol{\sigma}_a\cdot\mathbf{K})(\boldsymbol{\tau}_a\cdot\mathbf{A})e^{i\mathbf{K}\cdot\mathbf{r}_a}+(\boldsymbol{\sigma}_b\cdot\mathbf{K})(\boldsymbol{\tau}_b\cdot\mathbf{A})e^{i\mathbf{K}\cdot\mathbf{r}_b}], \quad (13)$$

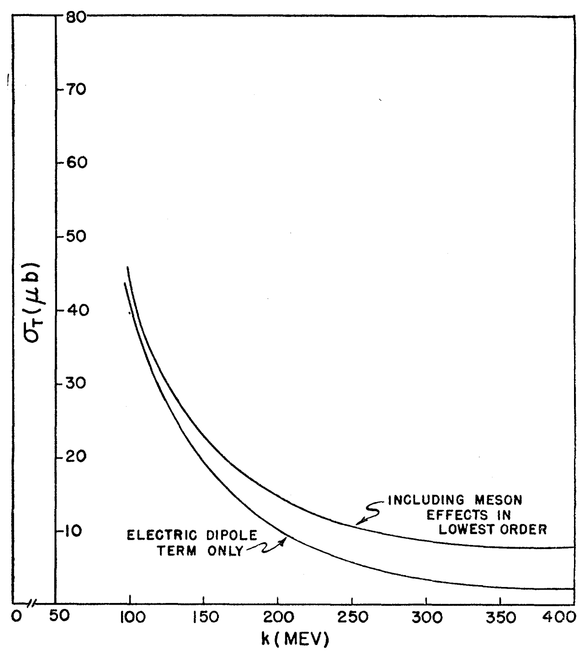


FIG. 2. Total cross section *vs* photon energy in the center-of-mass system, including only lowest-order effects.

[11] The notation here is as in Chew's article (reference 4). Higher multipoles in $H_n$ could be included, but have been found to be small by Marshall and Guth.

$$H_{nm}^{+}=(H_{nm}^{-})^{*},$$

$$H_3^{+}=i(4\pi)^{\frac{1}{2}}\frac{f}{\mu}eT\frac{1}{(2\omega)^{\frac{1}{2}}}\frac{1}{(2k)^{\frac{1}{2}}}[(\boldsymbol{\sigma}_a\cdot\hat{e})(\boldsymbol{\tau}_a\cdot\mathbf{a}^{\dagger})e^{-i\mathbf{K}\cdot\mathbf{r}_a}e^{i\mathbf{k}\cdot\mathbf{r}_a}+(\boldsymbol{\sigma}_b\cdot\hat{e})(\boldsymbol{\tau}_b\cdot\mathbf{a}^{\dagger})e^{-i\mathbf{K}\cdot\mathbf{r}_b}e^{i\mathbf{k}\cdot\mathbf{r}_b}],$$

$$H_m=-ie\mathbf{A}\cdot(\phi\boldsymbol{\nabla}\phi^{*}-\phi^{*}\boldsymbol{\nabla}\phi).$$

The nuclear current coupling can be written in a more convenient form by using Siegert's theorem[12]:

$$H_n=-ie(k/2)^{\frac{1}{2}}\hat{e}\cdot\boldsymbol{\nabla}-ie\cdot\mathbf{r}V_{\text{exchange}}-\sum_{a,b}\frac{(\mu_p+\mu_n)+(\mu_p-\mu_n)\tau_a{}^3}{2}\boldsymbol{\sigma}_a\cdot(\boldsymbol{\nabla}\times\mathbf{A}),\quad(14)$$

where $\mathbf{r}$ is the relative coordinate in the deuteron, and $V_{\text{exchange}}$ is the exchange part of the nuclear force, produced by the exchange of charged mesons. The extra term appearing here may be interpreted as being the contribution resulting from absorption of the photon by a virtual meson at very low energies. One would therefore expect this term to cancel, at low photon energies, the effects of the terms

$$H_{nm}^{-}\frac{1}{a}H_3+H_3\frac{1}{a}H_{nm}^{-}+H_{nm}^{-}\frac{1}{a}H_{nm}^{-}\frac{1}{a}H_m+H_{nm}^{-}\frac{1}{a}H_m\frac{1}{a}H_{nm}^{+}+H_m\frac{1}{a}H_{nm}^{+}\frac{1}{a}H_{nm}^{+},$$

and this in fact turns out to be the case. The part of $H_n$ remaining (that is, except for the $V_{\text{exchange}}$ term), is the ordinary electric and magnetic dipole photodisintegration, as computed by Marshall and Guth and Schiff.[2] At all but extremely low energies, the magnetic dipole term can be ignored, so that effectively

$$H_n=-ie(k/2)^{\frac{1}{2}}\hat{e}\cdot\mathbf{r},\quad(15)$$

and it must be remembered that the low-energy limit of the contribution coming from meson absorption of the photons must be subtracted from the remaining low-order terms.

The reasons for this choice of $H_n$, instead of the original one, are two. First, the present form allows a direct comparison with the earlier work which ignores explicit meson effects. Second, the form $\hat{e}\cdot\mathbf{r}$ emphasizes the deuteron wave function at large distances, where it is fairly well known, while the $\hat{e}\cdot\boldsymbol{\nabla}$ form emphasizes it at small distances, where it is not known.

The calculation of the lowest order terms is straightforward, using the approximations outlined above. The resulting cross section [to order $k/(p_f{}^2+u^2)^{\frac{1}{2}}$] is plotted in Fig. 2. The coupling constant is taken as $f^2=0.1$. The cross section ignoring meson effects (that is, taking the coupling to be $H_n$ only) is plotted on the same curve, for the same choice of deuteron wave function.

The angular distribution appearing here is symmetrical about 90°, being of the form $A+B\sin^2\theta+\sin^4\theta$.

As expected, no resonant effects appear, and the low-order terms serve merely as a fairly constant and quite small background at high energies. Comparison of the cross section coming from $H_n$ alone shows the explicit meson effects to be quite small.

An explicit calculation of the resonant terms in Eq. (10) requires a knowledge of the meson-nucleon scattering amplitude in the $(\frac{3}{2},\frac{3}{2})$ state off the energy shells. This has been taken as Gammel's approximate form,[13] which he found to agree very well with numerical solutions of the integral equation describing the meson-nucleon scattering. We thus take

$$(\mathbf{K}'\mathbf{a}'|R(a)|\mathbf{K}\mathbf{a})=\frac{1}{4\pi K'K}[3\mathbf{K}'\cdot\mathbf{K}-(\boldsymbol{\sigma}_a\cdot\mathbf{K}')(\boldsymbol{\sigma}_a\cdot\mathbf{K})]\tfrac{1}{3}[3\mathbf{a}'^{\dagger}\cdot\mathbf{a}-(\boldsymbol{\tau}_{\mathbf{a}}\cdot\mathbf{a}'^{\dagger})(\boldsymbol{\tau}_{\mathbf{a}}\cdot\mathbf{a})](K'|V_{\frac{3}{2}\frac{3}{2}}|K)\times\left\{1+\frac{\dfrac{1}{(K'|V_{\frac{3}{2}\frac{3}{2}}|K)}\displaystyle\int_0^{K_{\max}}\frac{K''^2dK''}{(2\pi)^3}\frac{(K'|V_{\frac{3}{2}\frac{3}{2}}|K'')(K''|V_{\frac{3}{2}\frac{3}{2}}|K)}{(k-w''+ie)}}{1-\displaystyle\int_0^{K_{\max}}\frac{K''^2dK''}{(2\pi)^3}\frac{(K''|V_{\frac{3}{2}\frac{3}{2}}|K'')}{k-w''+ie}}\right\},\quad(16)$$

where

$$(K'|V_{\frac{3}{2}\frac{3}{2}}|K)=4\pi\cdot\frac{8\pi f^2}{3\mu^2}\frac{K'K}{(w'w)^{\frac{1}{2}}}\frac{1}{k-w'-w+ie}.$$

We have chosen $f^2=0.1$, and a cutoff at the nucleon mass. These values, according to Gammel's tables, seem to fit the meson-nucleon scattering data reasonably well. It should be observed that the above form assumes a nonrecoiling nucleon. This means that in the photodisintegration, recoil is neglected except during the exchange of a meson between the two nucleons. Using this form for $R(a)$ and $R(b)$, the calculation of the cross section is again straightforward. The integral over the energy of the virtual meson which appears in the $H_3$ term has been evaluated numerically. The resulting total cross section, including the lowest-order terms, is indicated in Fig. 3. It is seen to be considerably too high for large photon energies, ranging from 20% to 50% above the experimental points. In view of the numerous approximations, it is questionable whether

[12] A. J. F. Siegert, Phys. Rev. 52, 787 (1937).

[13] J. L. Gammell, Phys. Rev. 95, 209 (1954).

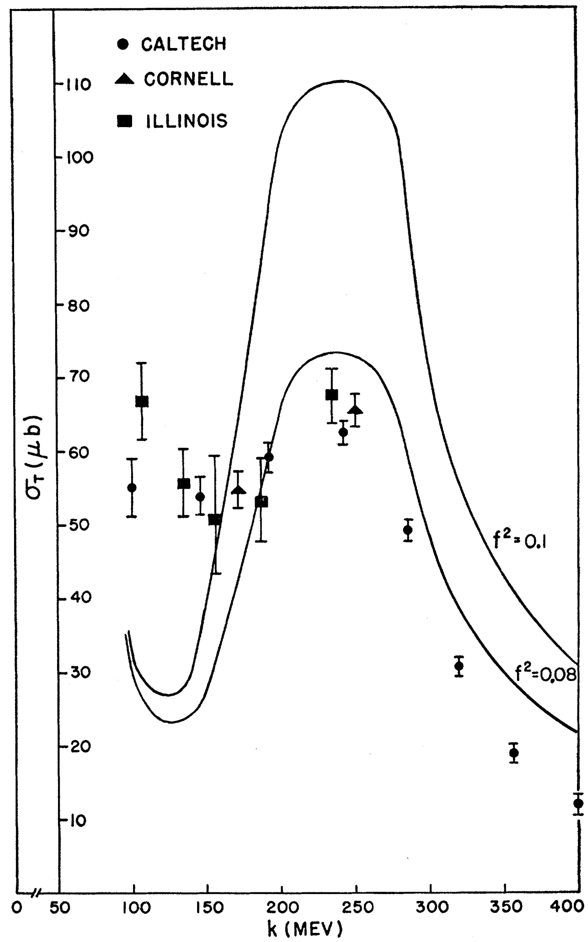


FIG. 3. Total cross section *vs* photon energy in the center-of-mass system, including resonant effects.

any better agreement than this could be anything but accidental. The theoretical results do, at least, predict the qualitative features of the resonance, such as its position and general shape.

It may be useful, however, to mention several points in the approximations which tend to overestimate the cross section. First, as was observed earlier, the choice of a different wave function for the deuteron could alter the cross section by up to 40% or so. A repulsive-core wave function would tend to reduce the results, by perhaps this amount. Second, the fact that multiple scattering in the deuteron has been neglected also tends to overestimate the result.[10] Finally, the coupling constant is not known completely accurately, and the resonant features of the cross section depend quite strongly on it. The choice of $f^2=0.1$ made here is, in fact, somewhat higher than the value generally accepted now. In order to indicate the sensitivity to the coupling constant, the cross section has also been computed using $f^2=0.08$, and making the reasonable assumption that $R(a)$ is not appreciably affected by this change. The resulting cross section is also shown in Fig. 3, and clearly agrees quite well with the experiments, except for the dip before the peak. The size of this dip is primarily due to a destructive interference between the perturbation and the three field terms. If the three-field term is omitted, and the larger coupling constant is used, the agreement is definitely improved. Unfortunately, there seems to be no valid reason for dropping the three-field term, as it is just about the same size as the anomalous moment term.

Regarding the angular distribution: an asymmetry is obtained from the interference of the two resonant terms of the form $\sin^2\theta \cos\theta$. This asymmetry becomes very small at low and high energies, and also gives the same cross section at 0° and 180°, which is apparently inconsistent with the experiments. It may be that a $\cos\theta$ dependence can only be found by investigating the $D$ state in the deuteron.

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