lead to the $j=\frac{1}{2}$ of the proton from a state of $J=\frac{3}{2}+$. Following the method of Hamilton, which permits the angular distribution to be expressed in terms of the most general quantum-mechanical features of angular momentum vectors, we take the incident γ-ray to define the z axis. The angular distribution of the reaction is then given by

$$W(\theta)=\sum_{m_i m_\pi m_\gamma}\sum\sum\sum|a(j_i l_\gamma; J; j_f l_\pi) \times(j_i, m_i; l_\gamma, m_\gamma|j_i, l_\gamma; J, m=m_i+m_\gamma)\mathfrak{D}_{l\gamma}{}^{m\gamma}(0) \times(j_f, l_\pi; J, m/j_f, m-m_\pi; l_\pi, m_\pi)Y_{l\pi}{}^{m}\ (\theta,\varphi)|^2, \quad (1)$$

where a is the matrix element (amplitude) for the reaction; $\mathfrak{D}$ and Y are the vector and ordinary spherical harmonics, respectively; $j_i=\frac{1}{2}$, m_i represents the initial proton state, $j_f=\frac{1}{2}$, m_f the final proton state, J, m the intermediate state angular momentum quantum numbers; and $(j_1, m_1, j_2, m_2/j_1, j_2; J, m=m_1+m_2)$ are the Clebsch-Gordon coefficients, tabulated by Condon and Shortley.[4] These calculations have been carried out for a number of possibilities and the results are given in Table I. Also given in the last column of Table I is the expected dependence of the cross section, or $|a|^2$, near threshold on the momentum, p, of the π-meson in the c.m. system (i.e., neglecting the dependence on γ-ray energy). An interesting feature of the above results is that $W(\theta)$ depends only on the values of J and l_γ. This is easily understood if one considers the inverse process; an incident π-meson along the z axis cannot alter the m-value of the system and therefore, irrespective of the value of l_π, leads always to the same intermediate states.

TABLE I. Angular and energy distributions in the photoproduction of π-mesons on nucleons.

γ-ray absorbed	Intermediate state	l of π-meson	$W(\theta)$	π-momentum dependence
Mag. dipole	1/2+	1	constant	p^3
Mag. dipole	3/2+	1	$2+3\sin^2\theta$	p^3
Elect. dipole	1/2−	0	constant	p
Elect. dipole	3/2−	2	$2+3\sin^2\theta$	p^5
Elect. quad.	3/2+	1	$1+\cos^2\theta$	p^3
Elect. quad.	5/2+	3	$1+6\cos^2\theta-5\cos^4\theta$	p^7

The ambiguity can, of course, be resolved by an observation of the energy dependence of the cross section near threshold. In the case of π^0 production, this turns out experimentally to be $\propto p^3$ of the meson (in the c.m. system) which, taken together with the angular distribution, indicates that the process is magnetic dipole absorption to a $J=\frac{3}{2}+$ intermediate state. In the case of π^+ production there is, however, strong indication of a first power p-dependence near threshold, which requires electric dipole absorption to a $\frac{1}{2}-$ intermediate state.

As can be seen from Table I, there are a variety of possible reactions, even if we confine ourselves to dipole-absorption processes. Furthermore, if more than one of these processes occurs in the same reaction, it is possible to have interference effects.[5] We confine our attention to two possibilities only: magnetic dipole absorption to $J=\frac{3}{2}+$, and electric dipole absorption to $J=\frac{1}{2}-$. The expected angular distribution is, then,

$$W(\theta)=|a|^2+2\,\mathrm{Re}(ab^*)\cos\theta+\tfrac{1}{2}|b|^2(2+3\sin^2\theta), \quad (2)$$

where b and a are, respectively, the amplitudes for the two processes.[6]

Now, the above analysis appears, at first sight, not to allow for the experimentally observed difference between π^0 and π^+ production on protons, since angular momentum considerations alone do not distinguish between the charge states of the meson and nucleon. Here, we can invoke the assumption of charge independence which introduces an additional quantum number into the system—the isotopic spin. The two reactions,

$$\gamma+p\rightarrow p+\pi^0$$

and

$$\gamma+p\rightarrow n+\pi^+,$$

differ only in the relative amplitudes of the final state wave function in the isotopic spin $\frac{3}{2}$ and $\frac{1}{2}$ components. In general, the reaction amplitudes are

$$a_J(\pi^0)=(\sqrt{\tfrac{2}{3}})a_{J,\,T=\frac{3}{2}}+(\sqrt{\tfrac{1}{3}})a_{J,\,T=\frac{1}{2}} \quad (3a)$$

for the first reaction, and

$$a_J(\pi^+)=(\sqrt{\tfrac{1}{3}})a_{J,\frac{3}{2}}-(\sqrt{\tfrac{2}{3}})a_{J,\frac{1}{2}} \quad (3b)$$

for the second. Following Brueckner and Watson,[2] we can choose $\sqrt{2}a_{\frac{1}{2},\frac{3}{2}}=-a_{\frac{1}{2},\frac{1}{2}}$. With this choice (made only on the basis of the experimental evidence on π^0 production), the dipole-absorption and the interference terms vanish for the π^0 process, but remain for the π^+ process. We are still left with two arbitrary amplitudes for the magnetic-dipole process, i.e., $a_{\frac{3}{2},\frac{3}{2}}$ and $a_{\frac{3}{2},\frac{1}{2}}$. There is, on the basis of the present photomeson experiments, no way of choosing these. However, the evidence from π-meson scattering on hydrogen has been interpreted[7] as indicating that $a_{\frac{3}{2},\frac{3}{2}}\gg a_{\frac{3}{2},\frac{1}{2}}$. The relative values of these two constants could also be obtained from a comparison of the cross sections for those parts of the π^0 and π^+ photoproduction reactions corresponding to magnetic-dipole absorption (the $2+3\sin^2\theta$ terms). In particular, for the assumption of $T=\frac{3}{2}$ production only, the π° cross section would be twice as great as the π^+ cross section.

While the evidence is as yet by no means conclusive, the results on π^0 production taken together with the relative magnitudes of the $\pi^\pm$ scattering processes are suggestive of a resonance corresponding to a proton "isobar" state of angular momentum $\frac{3}{2}$ in both ordinary and isotopic spin space.[2,7] However, the shape of the π^0 photoexcitation cross section is not well enough known to permit an evaluation of the "isobar resonance" constants. Further measurements, both on photoproduction and scattering, with improved accuracy and extended energy range, are required before an unambiguous answer can be given to the questions of the existence and properties of the isobar.

* This work was supported in part by the joint program of the U. S. Office of Naval Research and the U. S. Atomic Energy Commission.

[1] See preceding letter by Goldschmidt-Clermont, Osborne, and Scott [Phys. Rev. **89**, 329 (1953)] which also contains references to other experiments.

[2] K. A. Brueckner and K. M. Watson, Phys. Rev. **86**, 923 (1952).

[3] D. R. Hamilton, Phys. Rev. **58**, 122 (1940).

[4] E. U. Condon and G. H. Shortley, *The Theory of Atomic Spectra*, (Cambridge University Press, Cambridge, 1935), p. 76.

[5] $W(\theta)$ is now computed by an additional summation, in Eq. (1), taken inside the absolute value signs and over the possible intermediate J-values.

[6] It is also possible to include terms corresponding to magnetic dipole absorption to a $J=\frac{1}{2}+$ state, which gives rise to an interference term $\propto(3\cos^2\theta-1)/2$ which, however, does not lead to any asymmetry about 90° and which is not needed for the interpretation of the experiments now available.

[7] K. A. Brueckner, Phys. Rev. **86**, 106 (1952); Anderson, Fermi, Long, and Nagle, Phys. Rev. **85**, 936 (1952); Anderson, Fermi, Nagle, and Yodh, Phys. Rev. **86**, 793 (1952).

Infrared Photoconductivity Due to Neutral Impurities in Silicon

E. BURSTEIN, J. J. OBERLY, AND J. W. DAVISSON
Crystal Branch, Metallurgy Division, Naval Research Laboratory, Washington, D. C.

(Received October 6, 1952)

IMPURITY photoconductivity has been observed in silicon at liquid helium temperatures out to 38 microns. The regions of impurity photoconductivity and intrinsic photoconductivity are well separated so that it is possible to make an unambiguous interpretation of the data. The results, moreover, enable one to obtain information about the ionization energies of impurities and clearly demonstrate the value of optical measurements in the study of impurity levels in semiconductors.

Optical data obtained at liquid nitrogen temperature over the range 2 to 25 microns have previously demonstrated the absorption by neutral impurities in n- and p-type silicon involving the photoionization of bound charge carriers.[1] Efforts to detect

photoconductivity at liquid nitrogen temperatures due to this absorption were unsuccessful. An appreciable photoconductive response was observed, however, in various n- and p-type specimens at liquid helium temperatures. Using a residual ray monochromator with five crystal quartz plates and a rock salt shutter, it was possible to demonstrate a response as far out as 23 microns.[2] Spectral response measurements were subsequently carried out with a Perkin-Elmer monochromator having interchangeable NaCl, KBr, and KRS-5 prisms.

The spectral response of a relatively pure n-type silicon specimen containing 4×10^{15} charge carriers/cm^3 at room temperature is given in Fig. 1. The photoconductive response is proportional to

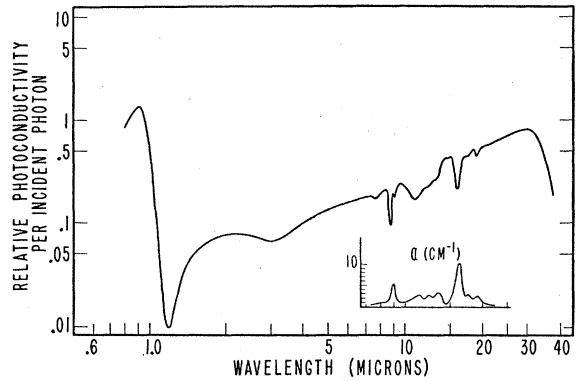

FIG. 1. The relative photoconductive response per incident photon of a relatively pure n-type silicon specimen containing 4×10^{15} charge carriers/cm^3 at room temperature. The strong dip in the photoconductive response at 1.2 microns is due to the superposition of a photonegative response at the absorption edge. The dips in the photoconductive response appearing between 8 and 24 microns correspond to the peaks in optical absorption due to lattice vibrations which compete with the photoionization absorption by the neutral impurities. The room temperature absorption spectrum of silicon due to lattice vibrations is given in the insert.

the light intensity, extends to 38 microns (the limit of measurement), and is a maximum at 32 microns. The time constant of the photoconductive response as determined by square wave chopping was found to be less than 10^{-4} second. The room temperature absorption spectrum of silicon due to lattice vibrations[3] is given in the insert in Fig. 1.

It is of interest to note that the dark resistance of the various n- and p-type silicon specimens at liquid helium temperatures was considerably smaller than the values expected from the ionization energies and concentrations of the impurities. By enclosing the specimen in a light-tight box with a shutter, it was found that the true dark resistivity of the relatively pure n-type silicon specimen at liquid helium temperature was greater than 10^{12} ohm-cm. With the shutter open to background radiation, the resistivity dropped to 5×10^{7} ohm-cm. The fractional change in resistivity of other specimens with higher impurity concentrations were, however, less than 1 percent when the shutter was opened and closed. In general, the larger the impurity content, the lower the residual resistance at liquid helium temperature, and also the smaller the photoconductive response. It appears therefore that specimens with relatively high impurity concentrations contains an appreciable fraction of impurities which have very small ionization energies.

According to the theory for the simple hydrogen model the photoionization absorption cross section of the neutral impurities may be expected to increase with wavelength to a maximum value at the ionization limit. Under conditions where the density of impurities is sufficiently small, such that the product of absorption constant and the thickness of the photoconductor is very much less than unity, the photoconductive response will be proportional to the absorption cross section and will likewise attain a maximum value at the ionization limit. The optical ionization energy of the neutral impurities may thus be obtained from the position of the peak in the spectral response curve, provided there is no appreciable distribution of impurity levels. The data of Fig. 1 yields a value of 0.04 ev for the optical ionization energy of the donor impurities in the relatively pure n-type specimen, which is somewhat smaller than the value 0.06 ev for the thermal ionization energy derived from Pearson and Bardeen's data for specimens of comparable impurity content.[4]

We wish to express our thanks to Dr. W. Shockley, Dr. J. Bardeen, and Mr. H. B. Briggs of the Bell Telephone Laboratories for helpful discussions and for supplying many of the specimens used in this investigation; to Dr. H. Stauss and Mr. J. Hino of the Metallurgy Division, Naval Research Laboratory, for making available specimens of pure silicon; and to Mrs. Bertha W. Henvis and Mr. H. Lipson of the Crystal Branch, Metallurgy Division, for their assistance in this investigation.

[1] Burstein, Oberly, Davisson, and Henvis, Phys. Rev. **82**, 764 (1951).
[2] Burstein, Oberly, and Davisson, Naval Research Laboratory Report of Progress, (July, 1950), p. 38 (unpublished).
[3] E. Burstein and J. J. Oberly, Phys. Rev. **78**, 642 (1950).
[4] G. L. Pearson and J. Bardeen, Phys. Rev. **75**, 865 (1949).

Asymmetric Fission

PETER FONG*
Department of Physics, University of Chicago, Chicago, Illinois
(Received October 2, 1952)

THE calculations by Frankel and Metropolis[1] have shown that at the saddle point of the energy in the fission process the nucleus is symmetric. However, at this point, there is practically no indication of a "neck," at which the deformed nucleus might break. The calculations by Hill[2] demonstrate that the fission process is slow enough that surface waves travel from one end to the other many times before a definite neck develops and fission occurs. In this letter it is proposed that the mode of fission is still undetermined at the saddle point. The concept of statistical equilibrium used by Bohr and Wheeler[3] is extended from the saddle point to a much later stage when the fission fragments are just about ready to come apart. Accordingly the number of quantum states at that later stage will determine the relative probability of different modes of fission. For convenience of calculation the situation at the breaking point is represented by a simplified picture, namely, that of the two fragment nuclei of mass A_1 and A_2 in contact.

The fact that the number of quantum states for asymmetric fission is larger than that for symmetric fission is due mainly to the fact that the internal excitation energy of the fragments at the breaking point is larger for the asymmetric than for the symmetric mode. According to the model adopted, the excitation energy at the breaking point is

$$E=M^*(A,Z)-M(A_1,Z_1)-M(A_2,Z_2)-E_{e1}-D. \qquad (1)$$

Here $M^*(A,Z)$ indicates the mass of the original excited fissioning nucleus, $M(A_1,Z_1)$ and $M(A_2,Z_2)$ the masses of the two fission fragments in their ground states, and E_{e1} the electrostatic repulsion between the two fragments. Since the nuclei are presumably highly deformed, a deformation energy D is introduced which decreases the energy available for distribution among the internal quantum states of the fragments. D is roughly independent of the mode of mass splitting and is estimated to be about 9 Mev. The term $-E_{e1}$, being proportional to Z_1Z_2, favors asymmetric fission. The mass terms, according to the liquid drop model mass values calculated by Metropolis and Reitwiesner,[4] favors symmetric fission. For example, the sum of masses of two equal fragments, Cd^{118}, is lower than that of two fragments of the experimentally most probable mode, Zr^{100} and Te^{136}, by 4.2 Mev. However, the liquid drop model mass values, compared to mass spectrometric ones, have shown errors of the order of 10 Mev. The masses of Cd^{118}, Zr^{100}, and Te^{136}, though not yet determined experimentally, can be extrapolated from the masses of stable nuclei[5] and from the parabolic dependence[6] of mass on charge number. The results show that the mass of $Cd^{118}+Cd^{118}$ is higher than