

rod. Temperatures reported here were deduced from the susceptibility of the salt (potassium chrome alum).⁶ The carbon resistors reached their highest resistance values approximately 25 minutes after the demagnetizing field was reduced to zero. All the data presented here are the result of transitions observed at times greater than 40 minutes after the demagnetization.

The superconducting transitions reported here were detected by observing the differential magnetic susceptibility $(\Delta M/\Delta H)_{H_a}$. The change in the magnetic moment (ΔM) of the sample is caused by the application of the incremental magnetic field ΔH produced by the primary winding of the mutual inductance. Field-induced transitions, such as shown in Fig. 1, as well as warm-ups in the presence of externally applied magnetic fields, yield the critical magnetic field data presented in Fig. 2. The magnetic fields employed here were always longitudinal to the mutual inductance coil system. Consequently, the appearance of the differential paramagnetic effect,⁷ such as depicted in Fig. 1, indicates that the specimen is exhibiting a good Meissner effect.

The data represent the results of two separate runs. Between the runs the Dewar was warmed to room temperature and the system was completely taken apart. The Ir specimens were taken off and reglued to a different section

of the copper rod. This placed the Ir samples in a secondary coil different from that used in the first run. Superconductivity was again detected for the Ir specimens and the results were consistent with those of the first run.

We are currently examining other high-purity metals to see if the occurrence of superconductivity is more widespread than previously believed and to see if the failure to become superconducting correlates with the presence of impurities with localized magnetic moment such as iron.

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¹Based on an extrapolation of the critical field data, corrected for the presence of a residual magnetic field of 1.5 gauss.

²B. T. Matthias *et al.*, Phys. Rev. Letters **5**, 542 (1960).

³A. M. Clogston *et al.*, Phys. Rev. **125**, 541 (1962).

⁴T. H. Geballe *et al.*, Phys. Rev. Letters **8**, 313 (1962).

⁵B. B. Goodman, Nature **167**, 111 (1951).

⁶Temperatures below 0.08°K are uncertain due to the unknown quality of thermal contact between the specimen and salt at these lower temperatures. In any event, they represent a lower limit on the transition temperatures.

⁷R. A. Hein and R. L. Falge, Jr., Phys. Rev. **123**, 407 (1961).

EXCITON-EXCITON INTERACTIONS AND PHOTOCONDUCTIVITY IN ORGANIC CRYSTALS

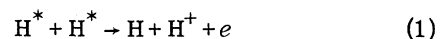
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Despite extensive study, the origin of photoconductivity in organic crystals remains obscure. Of the various suggestions which have been made we select for discussion that of Northrop and Simpson.¹ These investigators conclude from studies of photoconductance and fluorescence quenching in doped anthracene that exciton-exciton interactions are responsible for the promotion of electrons into conducting states. To render this mechanism plausible, it must be demonstrated that the transfer of excitation energy between two excited molecules with nonoverlapping electronic wave functions has a significant probability. In this Letter we consider the properties of a lattice of hydrogen atoms for which exact wave functions are available. It is found

that a transfer of excitation energy between degenerate states can occur very readily and therefore that exciton-exciton interaction can lead to photoconductivity. For an organic crystal, the breakdown of optical selection rules makes the process considered even more likely.

Consider first two hydrogen atoms a distance R apart, with R large enough that electron exchange may be neglected. Since the first excited state of H (and of many atoms and molecules) lies more than halfway to the ionization limit, only two quanta of excitation are required for ionization. From time-dependent perturbation theory the transition probability for the process



is, approximately,

$$T = \frac{4\pi}{k} \sum_{l,m} \left| \int \phi_1^A(\vec{r}_1) \phi_1^B(\vec{r}_2) \mathfrak{H} u_{klm}^A(\vec{r}_1) \phi_0^B(\vec{r}_2) d\tau_1 d\tau_2 \right|^2, \quad (2)$$

in atomic units. As usual, \mathfrak{H} is the interaction Hamiltonian, ϕ_i is a hydrogenic wave function for the first excited state if $i=1$ and for the ground state if $i=0$, and u_{klm} is a hydrogenic continuum wave function. When only dipole-dipole interactions are considered, then

$$T \approx 2.09/R^6 \quad (3)$$

in atomic units (au).

Consider next a simple tetragonal lattice with N hydrogen atoms per cubic centimeter and unit cell characterized by the basis vectors $R\hat{i}$, $R\hat{j}$, $R'\hat{k}$. We describe two excitons in the crystal by the wave function

$$\begin{aligned} \psi_{\vec{k}, \vec{k}'}(\lambda, \lambda') \\ = \frac{1}{[N(N-1)]^{1/2}} \sum_{\vec{I} \neq \vec{I}'} \sum_{\vec{I}'} e^{i\vec{k} \cdot \vec{I} + i\vec{k}' \cdot \vec{I}'} |\lambda, \lambda', \vec{I}, \vec{I}'; \vec{r}_1 \cdots \vec{r}_N\rangle, \end{aligned} \quad (4)$$

where \vec{I} is a lattice vector,

$$\begin{aligned} |\lambda, \lambda', \vec{I}, \vec{I}'; \vec{r}_1 \cdots \vec{r}_N\rangle \\ = \phi_{\lambda}^{\vec{I}}(\vec{r}_l) \phi_{\lambda'}^{\vec{I}'}(\vec{r}_{l'}) \prod_{i=1}^N \phi_0^{(i)}(\vec{r}_i), \end{aligned} \quad (5)$$

$i \neq l, l'$

and ϕ_{λ} is a $2p_x$, $2p_y$, or $2p_z$ electronic wave function. Similarly, the electron in the continuum state is represented by

$$\begin{aligned} \Psi_{klm}(\vec{p}) = \frac{1}{N^{1/2}} \sum_{\vec{L}} e^{i\vec{p} \cdot \vec{L}} |klm; \vec{L}\rangle, \\ |klm; \vec{L}\rangle \equiv u_{klm}^{\vec{L}}(\vec{r}_L) \prod_{i=1}^N \phi_0^{(i)}(\vec{r}_i). \end{aligned} \quad (6)$$

$i \neq L$

To evaluate the transition probability between the states (4) and (6), it will be assumed that (1) dipole-dipole interactions alone need be considered, (2) only nearest neighbor interactions are of importance, and (3) the value of k which satisfies conservation of energy is obtained from $2\epsilon_1 = \epsilon_0 + \frac{1}{2}k^2$, where ϵ_0 and ϵ_1 are the energies of the ground state and first excited state of an atom,

respectively. Then the probability of a transition to the continuum states is

$$\begin{aligned} T_x &= \frac{8.56}{NR^6} [-2 \cos k_{rx} R \cos \frac{1}{2} K_{cx} R + \cos k_{ry} R \cos \frac{1}{2} K_{cy} R \\ &\quad + (R/R')^3 \cos k_{rz} R' \cos \frac{1}{2} K_{cz} R']^2, \\ T_y &= \frac{8.56}{NR^6} [-2 \cos k_{ry} R \cos \frac{1}{2} K_{cy} R + \cos k_{rx} R \cos \frac{1}{2} K_{cx} R \\ &\quad + (R/R')^3 \cos k_{rz} R' \cos \frac{1}{2} K_{cz} R']^2, \\ T_z &= \frac{8.56}{NR^6} [\cos k_{rx} R \cos \frac{1}{2} K_{cx} R + \cos k_{ry} R \cos \frac{1}{2} K_{cy} R \\ &\quad - 2(R/R')^3 \cos k_{rz} R' \cos \frac{1}{2} K_{cz} R']^2, \\ \vec{k}_r &\equiv \frac{1}{2}(\vec{k} - \vec{k}'), \\ \vec{K}_c &\equiv (\vec{k} + \vec{k}'), \end{aligned} \quad (7)$$

for $2p_x$, $2p_y$, and $2p_z$ excitons. Under ordinary conditions $\vec{k}_r \cdot \vec{R}$ and $\vec{K}_c \cdot \vec{R}$ are small relative to unity and Eqs. (7) may be expanded in power series form. For the case that $R' \gg R$, and with the n excitons per cubic centimeter randomly distributed, the rate of production of ionized atoms is

$$\mathfrak{R}_{\text{au}} \approx \frac{1}{3} n^2 (8.56/NR^6), \quad (8)$$

in atomic units. In the more ordinary units of ionizations per second,

$$\mathfrak{R} \approx 11.6 \times (n^2/NR^6) \times 10^{16}. \quad (9)$$

As a crude approximation to an organic crystal for order-of-magnitude estimates, we take $n = 3.5 \times 10^{12}$, $N = 6 \times 10^{21}$, $R = 20$ au, $R' = 100$ au, and therefore

$$\mathfrak{R} \approx 3.7 \times 10^{12} \text{ ionizations/sec.}$$

Northrop and Simpson¹ estimated the rate of production of charge carriers in anthracene to be $1.2 \times 10^{17} \text{ sec}^{-1}$ and reached the conclusion that the mobility of the positive charge carrier has to be $5 \times 10^{-3} \text{ cm}^{-2} \text{ V}^{-1} \text{ sec}^{-1}$. Their value for the rate of production, when corrected to yield the result of recent measurements of the

mobility,^{2,3} is $3.5 \times 10^{12} \text{ cm}^{-3} \text{ sec}^{-1}$ in agreement with the theoretical estimate quoted. Because this calculation should not be directly applied to an organic crystal, and because Northrop and Simpson's determination of the rate of carrier production is subject to considerable uncertainty, our result should be regarded as a demonstration of the plausibility of the mechanism whereby charge carriers are produced from exciton-exciton interactions in a molecular crystal. Further details of this investigation and the development of techniques suitable to the proper study

of organic crystals will be presented in a separate publication.

We wish to thank the National Science Foundation and the Air Force Office of Scientific Research (Grant 61-52) for financial support.

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SPIN STRUCTURE OF THE MESON-NUCLEON AND NUCLEON-NUCLEON SCATTERING AMPLITUDES AT HIGH ENERGIES

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(Received April 23, 1962)

It has been recently shown¹ that generalized unitarity relations obtained by one of us (V.G.)² lead to the existence of relations between the amplitudes of the various processes at high energies when asymptotically these processes are caused by a single Regge pole.^{3,4}

In this note we shall consider the vacuum pole⁴ in meson-nucleon and nucleon-nucleon scattering amplitudes in an attempt to analyze their spin structure and the relation between them.

Let us write a general expression for the amplitudes of meson-nucleon scattering G and nucleon-nucleon scattering H^5 :

$$G = a + b\hat{q}, \quad q = \frac{1}{2}(k_1 + k_1'), \quad (1)$$

$$H = H_1(s, t) + H_2(s, t)(\gamma^{(1)}p + \gamma^{(2)}k) + H_3(s, t)(\gamma^{(1)}p)(\gamma^{(2)}k) + H_4(s, t)(\gamma_5^{(1)}\gamma^{(1)}p)(\gamma_5^{(2)}\gamma^{(2)}k) + H_5(s, t)\gamma_5^{(1)}\gamma_5^{(2)},$$

$$p = \frac{1}{2}(p_2 + p_2'), \quad k = \frac{1}{2}(p_1 + p_1'). \quad (2)$$

Here $k_1, k_1'; p_1, p_2; p_1', p_2'$ are initial and final pion and nucleon momenta, respectively. As far as the vacuum pole, which is responsible for the amplitudes at high energies and low t , is of interest to us, these amplitudes are unit matrices in isobaric space. It can be seen that terms containing $\gamma_5[H_4(s, t), H_5(s, t)]$ cannot be generated by

the vacuum pole and are therefore asymptotically small. Thus at high s and low t , the nucleon-nucleon amplitude contains only three invariant functions. As the vacuum pole appears only in the part of the amplitude symmetric under permutation of s and u ,⁴ the same expressions describe the scattering of antiparticles by nucleons. As in the case of spinless particles, let us consider the partial-wave expansion of amplitudes in the annihilation channel.

For the meson-meson scattering amplitude we have the usual expression

$$f(s, t) = \sum_j (2j+1) f_j(t) P_j(Z_f). \quad (3)$$

To expand two-meson annihilation and nucleon-nucleon scattering amplitudes, we shall use the helical amplitudes⁶

$$\langle |G| \lambda', \lambda \rangle = \sum_j (2j+1) g_j(\lambda', \lambda) d_{\lambda-\lambda', 0}^j(Z_g),$$

$$\langle \lambda_2', \lambda_2 | H | \lambda_1', \lambda_1 \rangle$$

$$= \sum_j (2j+1) \langle \lambda_2', \lambda_2 | H^j | \lambda_1', \lambda_1 \rangle d_{\lambda_1-\lambda_1', \lambda_2-\lambda_2'}^j(Z_h), \quad (4)$$

Z_f, Z_g, Z_h being cosines of scattering angles in