$$V(r) = -\frac{V_0}{1+\exp\left(\frac{r-R}{a}\right)} - \lambda V_0\left(\frac{\hbar}{m_\pi c}\right)^2 \frac{1}{r}\frac{d}{dr}\left[1+\exp\frac{(r-R)}{a}\right]^{-1}(\mathbf{l}\cdot\mathbf{s});$$

$$R = r_0 A^{1/3}.$$

The parameters were chosen from data on the elastic scattering of electrons and on the energy of detachment of the neutron from the product nucleus [5]. The results are listed in the table and are close to the experimental values.

In Fig. 2, the function $F_{1,1,1/2}(p)$ for ^{13}C is compared with the corresponding harmonic-oscillator function.

In spite of the fact that the radial wave functions differ negligibly within the limits of the nuclear radius, the momentum distribution at large momenta (≥ 2 F^{-1}) are entirely different in magnitude and in form (they practically coincide at small momenta). The function $F_{1,1,1/2}(p)$ in the Saxon-Woods potential at large momenta experiences oscillations that are possibly connected with reflections from the "sharp" boundary of the potential.

Figure 3 shows the angular distribution of the reaction $p + {}^{12}C \to \pi^+ + {}^{13}C$ at a proton energy 600 MeV. A characteristic feature is the presence of minima which are missing (at the considered proton energy) from the angular distribution obtained when p-n correlations are taken into account.

The present results show that the available experimental data can be described within the framework of the single-particle model of the nucleus. Therefore the contribution of the p-n correlations to the reactions in question is possibly smaller than that obtained in [2], and for their study it is necessary to have more detailed and exact experimental data.

One of the authors (A. Gridnev) is grateful to S.P. Kruglov for interest in the work.

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CONCERNING ONE MECHANISM OF THE PHOTOCONDUCTIVITY OF DISORDERED SEMICONDUCTORS

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Submitted 6 June 1972
ZhETF Pis. Red. 16, No. 2, 109 - 111 (20 July 1972)

As already noted by R. Kuyper and one of us, the hopping conductivity σ can be due to energy exchange between the carriers and any other elementary

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excitations, not only phonons. In particular, it is natural to investigate the hops stimulated by photons. The smallness of the coupling constant (in this case σ is proportional to the square of the fine-structure constant) seemingly makes this process inconvenient. However, there are two advantages here compared with hops stimulated by phonons. First, the photon energy is not limited by the Debye temperature. Consequently, there are no energy limitations on single-phonon processes, and long-distance hops are no longer obligatory. Second, in the presence of suitable added illumination, the number of photons of the required frequency turns out to be independent of the temperature. Accordingly the exponential factor (Boltzmann or Mott) is eliminated from σ. We arrive at a specific type of conductivity in which the radiation energy is consumed in stimulation of the hops[2]. Obviously, this phenomenon can be observed if the spectral composition of the additional illumination eliminates the possibility of photoconductivity of the usual type. At low temperatures the latter denotes that the frequency of light ω should satisfy the conditions $\hbar\omega < E_c - F$ and $\hbar\omega < F - E_v$, where E_c and E_v are the limits of the forbidden band, understood as the gap for the mobility, and F is the Fermi level (which lies by assumption sufficiently deep in the forbidden band).

A detailed calculation of the photon-stimulated hopping conductivity σ_{phot} is just as difficult as in the case of phonon processes (σ_{phon}), and the result depends both on the solution of the calculation problem and on the statistics of the random discrete levels. One can, however, estimate the lower limit of the ratio $\sigma_{phot}/\sigma_{phon} = \eta$, assuming that the indicated factors influence in approximately the same manner both σ_{phon} and σ_{phot} (actually percolation can only enhance the role of σ_{phot}).

Obviously, it suffices to consider the case of monochromatic illumination. It is necessary here to distinguish two cases:

$$\text{a)}\ |\,|F| - \hbar\omega\,| \sim |F|\ (\omega \equiv \bar{\omega}),\quad \text{b)}\ |\,|F| - \hbar\omega\,| \ll |F|\ (\omega \equiv \omega).$$

In the first case, the electrons are excited to levels that are sufficiently close to F (as in the case of interaction with phonons), and in the second they are excited to higher levels close to E_c: in this case σ_{phot} differs from the corresponding value in the case (a) by the factor

$$\xi \approx 10^3 \left(\frac{F}{\hbar\omega}\right)^2 \frac{\rho(F+\hbar\omega)}{\rho(F)} \left|\frac{F+n\omega}{F}\right|^{3/2} \frac{n^3(\omega)}{n^3(\bar{\omega})}\,\frac{\omega}{\bar{\omega}}\,. \qquad (1)$$

The factor 10^3 comes here from the integrals with respect to the coordinates, which contain powers of the distance between centers R and the overlap factors $\exp[-(\gamma_1 - \gamma_2)R]$, where γ_1 and γ_2 are the reciprocal radii of electron localization in states between which the hopping takes place. $\gamma_1 \approx \gamma_2 \approx \gamma(F)$ in case (a) and $\gamma_2 R \lesssim 1$ in case (b). When account is taken of the factor R^{10}, which appears when the corresponding matrix element is calculated with asymptotic wave functions of the discrete spectrum, this introduces into ξ a factor $2^{11} \approx 10^3$. We note that when $\gamma R \lesssim 1$ the use of the asymptotic form is strictly speaking no longer justified; it is easily seen, however, that allowance for the additional power-law dependence of the wave function on the coordinates only overestimates the value of ξ. Further, ρ is the density of states; the absolute-value sign is used only because we choose E_c as the zero point of the energy;

[2] An analogous "acoustic conductivity" process, namely hopping electric conductivity stimulated not by thermal but by sound waves introduced from the outside, would be of interest.

$n(\omega)$ is the refractive index at the frequency ω. The right-hand side of (1) has a maximum at $\hbar\omega \simeq |F| - W_1$, where the characteristic energy W_1 ($W_1 \leq |F|$) is the solution of the equation $d \ln \rho/dW = 3W/2$.

Wishing to obtain a lower limit for ξ, we compare σ_{phot} (in case (a)) with the hopping conductivity stimulated by single-phonon transitions; for concreteness we have in mind acoustic phonons and the simplest variant of the deformation-potential method. We confine ourselves here to a region of not too high temperatures, when [1 - 3] $\sigma_{phon} \sim \exp[-(T_0/T)^{1/4}]$, where T is the absolute temperature and T_0 is a constant.

We introduce the following notation: E_1 is the deformation potential, $\omega_m = T_m/\hbar$ is the Debye frequency, s is the speed of sound, d is the density of the material, I is the energy flux in the light wave, and m_0 is the mass of the free electron. We introduce also the characteristic energy values W_2, the "effective mass" of the electron m, the "atomic mass" M, and the energy flux I_0, assuming that

$$\int_{E_v}^{F} \rho(W)\,dW = \rho(F)W_2, \quad \frac{\hbar^2\gamma^2(F)}{2m} = |F|, \tag{2}$$

$$M = \frac{ds^3}{\omega_m^3}, \quad I_0 = \frac{T_m^4}{2\pi^2\hbar^3 s^2} \tag{3}$$

We then obtain for η in the case (a)

$$\eta = \frac{e^2 n^3(\bar{\omega})}{hc} \frac{IW_2|F|^3\bar{\omega}\, mMs^2}{I_0 k^2 T T_0 E_1^2 \omega_m m_0^2 c^2} \exp\left\{\left(\frac{T_0}{T}\right)^{1/4}\right\}. \tag{4}$$

The increase of η with increasing frequency of the light is due to the fact that we calculated σ_{phot} by using the dipole approximation (this is apparently justified), and the characteristic photon frequency was assumed equal to ω_m (this underestimates η). The value of η in case (b) is obtained from (4) by multiplying by ξ.

Obviously, the considered type of photoconductivity can be of experimental interest when $I > I_{cr}$, where the critical value I_{cr} is determined from the condition $\eta = 1$. Assuming for estimating purposes $n(\omega) = 7$, $m = m_0$, $|F| = 0.3$ eV, $E_1 = 5$ eV, $kT_m = 10^{-2}$ eV, $T = 81°K$, $M = 3 \times 10^5 m_0$, $s = 5 \times 10^5$ cm/sec, $T_0 = 10^8 °K$, and $W_2 = 0.1$ eV, we see that $I_0 \simeq 0.8 \times 10^5$ W/cm^2. Then, according to (4), $I_{cr} = 0.45\omega_m/\omega$ W/cm^2.

In the case (b) this value is decreased by a factor ξ. Choosing the frequency in optimal fashion and assuming as an estimate $\rho(-W_1)|_{\rho(F)} = 5$, we obtain in case (b) $I_{cr} \simeq 5 \times 10^5(\omega_m/\omega) \simeq 2.5 \times 10^{-6}$ W/cm^2.

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PLASMA HEATING IN CLOSED MAGNETIC TRAPS BY INJECTION OF FAST NEUTRAL ATOMS

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Submitted 7 June 1972
ZhETF Pis. Red. 16, No. 2, 112 - 114 (20 July 1972)

One of the methods of heating plasma ions in modern closed magnetic traps is by injecting into the plasma fast hydrogen atoms with energies of several tens or hundreds of keV [1, 2, 12, 13]. The fast neutral atoms are produced by neutralization of an ion beam in a charge-exchange gas target. When these neutral atoms enter the plasma they are ionized, captured in a trap, and transfer energy to the plasma particles, which maintain this energy during a characteristic time τ_e. It is shown in the present article that replacement of the hydrogen atoms by lithium atoms having the same energy makes this plasma heating method much more effective and preferable for use in existing experimental setups and in those under construction. Let us consider the main factors that determine the effectiveness of this method of plasma heating.

1. The rate of energy transfer from a fast ion with energy W to ions and electrons of a plasma is given by the relations [3, 4][1)]

$$\frac{dW}{dt_{ii}} = -1.8\cdot10^{-7}\frac{A_1^{1/2}Z^2nL}{W^{1/2}A_2}\ \text{eV/sec}, \tag{1}$$

$$\frac{dW}{dt_{ie}} = -3.2\cdot10^{-9}\frac{Z^2nWL}{A_1T_e^{3/2}}\ \text{eV/sec},^{2)} \tag{2}$$

where A_1 and Z are the mass number and charge of the fast ion, A_Z is the mass number of the plasma ions, L is the Coulomb logarithm, M and m are the masses of the proton and of the electron, n is the plasma concentration in cm^{-3}, and T_e is the electron temperature (we assume below that $A_2 = 1$ and $L = 15$ [3]). The limiting energy at which $dW/dt_{ii} = dW/dt_{ie}$ is [4]

$$W_2 \approx 16T_eA_1\ . \tag{3}$$

When $W < W_2$, the fast ions transfer energy predominantly to the plasma ions, and when $W > W_2$ to the electrons. The time of cooling of the fast ion to $W \approx T_j$, T_e is [12]

$$\tau = -\int_0^W \frac{dW}{dW/dt} \approx 2.6\cdot10^7\frac{A_1T_e^{3/2}}{Z^2n}\ln[1 + (W/W_2)^{3/2}] \tag{4}$$

It is necessary to substitute in (1) - (4) $A_1 = Z = 1$ for protons, and $A_1 = 7$ and $Z = 3$ for lithium ions (it is easy to see that the time that the lithium

[1)] The relations used in [2] are in error, since they determine the relaxation time of the translational momentum of the fast ion, and not the energy relaxation time.

[2)] $W < A_1MT_e/m$.