At W_{des} there is produced in AHC a dense grid of interacting electrons and holes with spacing 30 - 50 Å. We assume that their condensation leads to the formation in the AHC of macroscopic regions with an excitation-energy density greatly exceeding the average value. Keldysh [6] was the first to consider the occurrence of "electron-hole drops" in semiconductors, where the exciton has a large radius, a ∿ 100 A, and a low binding energy, ε ∿ 0.02 eV. Therefore the electrons and holes combine into one Wannier-Mott-Keldysh condenson with equilibrium particle density 10^{17} - 10^{18} cm^3, which are stable at low temperature. In AHC, a is of the order of the interatomic distance and ε ∿ 1 eV. We therefore assume that the system of electrons and holes at room temperature is capable of contracting into Frankel condensons with equilibrium density 10^{21} - 10^{22} cm^{-3}. W can then reach the binding-energy density, which is certainly sufficient for destruction.

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BEHAVIOR OF LARGE-CURRENT ELECTRON BEAM IN A DENSE GAS

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It was observed in [1] that a 50-kA beam of 3-MeV electrons traverses a distance of 30 cm in air at atmospheric pressure before it spreads apart. This is larger than the mean free path of one electron by more than one order of magnitude.

We have investigated this effect in different gases in the pressure interval from 10^{-2} Torr to 1.6 atm. A beam of electrons from the "Neptune" accelerator passed through the anode foil into a dielectric chamber of 18 cm diameter and 80 cm length. The bulk of the experiments were performed at an electron energy 660 keV and a beam current 12 kA and duration 40 sec. The beam parameters in the drift chamber were measured with a calorimeter that could be moved along the axis. At the same time, we photographed the glow of the plasma produced by the beam. Figure 1 shows photographs in air and in helium. Figure 2 shows the energy distribution, along the chamber, of an electron beam incident on the calorimeter with a diameter equal to the chamber diameter. A comparison of these curves shows that the plasma glow intensity agrees with the calorimetric measurements. We see that in air at atmospheric pressure the beam breaks up at a length ∿12 cm. The penetration length L of the beam in air decreases from 20 to 10 cm when the pressure is changed from 0.4 to 1.6 atm. At lower pressures, L increases rapidly. In helium we were unable to measure L, since it turned out to be larger than the length of the drift chamber in the pressure interval up to 1.6 atm.

These phenomena are apparently connected with the appearance of a unique beam instability in a strongly dissipative medium. Under the conditions

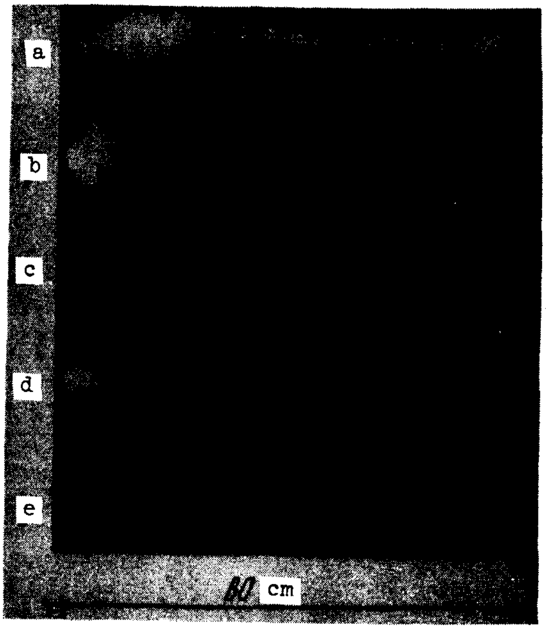

Fig. 1

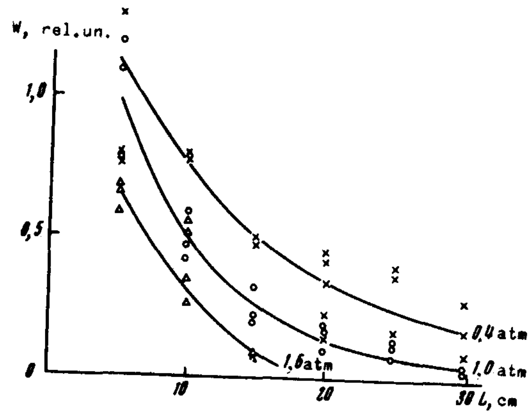

Fig. 2

Fig. 1. Glow of plasma produced by an electron beam: a) air, P = 10^{-2} atm, b) air, P = 0.4 atm, c) air, P = 1.0 atm, d) air, P = 1.6 atm, e) air, P = 1.6 atm.

Fig. 2. Energy distribution, along the chamber axis, of the beam incident on the calorimeter: a) P = 0.4 atm, b) P = 1.0 atm, c) P = 1.6 atm. The gas is air.

of the present experiment, the electrons of the relatively dense plasma ($n = 10^{14} - 10^{16}$ cm^{-3}) produced by ionization of the gas by the beam electrons has a high scattering frequency ($\nu_e \sim 4 \times 10^{12}$ sec^{-1} at 1 atm). Under these conditions, the dispersion equation describing the behavior of small perturbations of the beam density is

$$1 = \frac{\omega_p^2}{i\omega\nu_e} + \frac{\omega_p^2}{(\omega - k_n u)^2}\,\frac{n'}{n}\left(\frac{1}{\gamma}\,\frac{k_\perp^2}{k^2} + \frac{1}{\gamma^3}\,\frac{k_\parallel^2}{k^2}\right). \qquad (1)$$

Here n' and n are the particle concentrations in the beam and in the plasma, u is the beam velocity, $\gamma = (1 - u^2/c^2)^{-1/2}$, and $k_\parallel$ and $k_\perp$ are the components of the wave vector of the perturbation along and across the beam propagation direction. This equation has an unstable solution that leads to bunching of the beam [2]

$$\omega = k_\parallel u + \delta, \quad I_m\delta = \frac{1}{\sqrt{2}}\left[k u \nu_e \frac{n'}{n}\left(\frac{1}{\gamma}\,\frac{k_\perp^2}{k^2} + \frac{1}{\gamma^3}\,\frac{k_\parallel^2}{k^2}\right)\right]^{1/2}. \qquad (2)$$

For a radial mode having no zeroes on the beam axis, instability exists if

$$I_m\delta > k_\parallel \overline{\Delta u}, \qquad (3)$$

where Δu is the thermal scatter of the beam velocities along the propagation direction. Physically this condition means that the increase of the perturbation should be faster than its decay due to the velocity scatter. From (2) and (3) we can estimate the maximum increment, the wave vector of the most unstable perturbations, and the length over which the perturbation decreases by a factor e

$$(I_m\delta)_{max} \approx \frac{n'}{n\gamma}\nu_e\frac{u}{\Delta U}, \quad k_\parallel \approx \frac{n'}{n\gamma}\ \frac{\nu_e}{u}\left|\frac{u}{\overline{\Delta u}}\right|^2, \quad k_\parallel \leqslant k_\perp \leqslant \frac{1}{r},$$

$$\frac{u}{(I_m\delta)_{max}} \approx \gamma\frac{n}{n'}\ \frac{\overline{\Delta u}}{\nu_e}. \tag{4}$$

Bunching of the beam leads to its deceleration as a result of dissipation of the alternating current $j \simeq e\omega\delta n/k$ excited in the plasma. This current is produced by an electric field $E = e\omega\delta n/k\sigma$, and the dissipation of the beam energy $n\varepsilon$ in the plasma is described by the relation

$$\frac{(e\omega\delta n)^2}{k^2\sigma} \approx \frac{d}{dz}n'\varepsilon u. \tag{5}$$

For an already developed instability, we can put in this estimate $\delta n \sim n'$. In addition to the deceleration effect, the decrease of the energy flux can be connected with the drift of the beam particles to the walls as a result of the increase of the radial velocities under the influence of the electric field of the oscillations. Therefore the distance L traversed by the beam should lie between the values

$$\gamma\frac{n}{n'}\ \frac{\overline{\Delta u}}{\nu} \leqslant L \leqslant \gamma\frac{n}{n'}\ \frac{u}{\nu_e}. \tag{6}$$

The plasma concentration is due to gas ionization by the relativistic electrons and also by the plasma electrons in the region where the beam slows down. This secondary ionization leads, at sufficiently low pressures, to gas breakdown. If there is no breakdown, then the concentration of the electrons is determined from the relation $n_e = (n'n_0\langle\sigma v\rangle_i)^{1/2}/\alpha$ (α is the recombination coefficient), under the condition that the characteristic recombination time $1/\alpha n_e$ is smaller than the beam duration τ_p. On the basis of [3, 4] we have $\alpha = 10^{-7}$ cm³/sec for air at $T_e \sim 1$ eV. Then $1/\alpha n_e \lesssim \tau_p$ at $P > 0.1$ atm. Therefore, in accordance with (6) we have

$$\gamma\frac{\Delta U}{\langle\sigma V\rangle_e}\sqrt{\frac{\langle\sigma V\rangle_i}{n'n_0\alpha}} < L. \tag{7}$$

Let us estimate the value of L for an air pressure of 1 atm. At $n' = 2 \times 10^{11}$ cm^{-3}, $n_e \simeq 2 \times 10^{15}$ cm^{-3} and $\nu \simeq 4 \times 10^{12}$ sec^{-1} we have $L > 90(\Delta m/c)$ cm. The value of L is close to that observed in experiments in air, if we put $\Delta U \simeq 5 \times 10^9$ cm/sec. The dependence of the length L on the pressure agrees with our experimental results at $P > 0.4$ atm.

At $P < 0.4$ atm in air, L increases more rapidly than shown by formula (7). This circumstance can be attributed to the additional ionization during the gas breakdown. The electric field resulting from the instability is larger the smaller the length of the instability development. Therefore in helium, where the breakdown field is smaller than in air by a factor 4 - 5 [5], the beam penetration length is larger. We note that the dielectric strength of the gas may be decreased by its heating by the powerful beam, which leads to an increase in the propagation length.

Formula (7) agrees also with measurements performed earlier by Link [1]. In these experiments, in air at atmospheric pressure, L is approximately 4 times larger than in our case, owing to the increase of γ and of the energy in the beam.

In conclusion, we note once more that a substantial fraction of the energy can be converted into heat in the phenomenon in question. It can therefore be used for a collective heating of a superdense plasma.

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SEPARABLE EXPANSION OF THREE-PARTICLE AMPLITUDES

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In the present article, using the simplest separable interaction as an example, we consider certain properties of the Hilbert-Schmidt (HS) expansion [1 - 3] for amplitudes in the three-body problem [4]. An analysis of an expansion of this type is very important both from the point of view of clarifying the analytic structure of the three-particle amplitudes, and in connection with the possibility of using the HS expansion to solve integral equations for four particles [5, 6].

Our result reduces to the following. For identical particles we can introduce two three-particle amplitudes X(z) and Y(z) [7] (z is the energy parameter), the first of which is connected with the amplitude of physical processes (elastic n-d scattering, disintegration of a deuteron) in a system of three particles. Since the equations for the amplitudes X(z) and Y(z) separate, it follows that Y(z) does not correspond to any physical quantity and is not used in practice in three-particle calculations. The amplitudes X(z) and Y(z) are in the general case ($N \times N$) matrices, where N is the number of terms in the separable expansion of the two-particle T-matrix. Symbolically, the equation for X(z) can be written in the form

$$X(z) = V(z) - V(z)\Delta(z)X(z), \tag{1}$$

where V(z) is the effective three-particle potential, and $\Delta(z)$ is a diagonal matrix describing the propagation of a free particle and of a pair of interacting ing particles. The equation for Y(z) differs from (1) in the substitution $V \to V/2$ and in the reversal of the sign of the second term. It follows from (1) that the HS expansion of the amplitude X(z) coincides formally with the HS expansion of the two-particle amplitude. Each term of this expansion contains