One more remark. It is easy to verify that the relative motion of vortices is a Liouville motion, so that it is possible to write down a microcanonical distribution with respect to two invariants of the motion (1.2) and (1.5). The decisive parameter of this distribution in the case of identical vortices is the quantity $\theta$ introduced in Ref. 4 (the analog of the temperature). It is of interest to investigate in detail in the future the stochastic properties of the system of four and more vortices so as to verify, for example, the applicability of the microcanonical distribution in a definite range of values of $\theta$ (the calculations performed in the present paper correspond to the case $\theta \sim 2.4$).

We are grateful to V. I. Arnol'd and Ya. G. Sinaĭ for a useful discussion of the results and to V. E. Zakharov for a stimulating discussion during one of the stages of the work.

1) A graph is called planar if it can be drawn on a plane (or on a sphere) without self-intersections.

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# Recombination decay of cryogenic helium plasma

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An experimental and theoretical study is reported of the kinetics of electron density ($n = 10^{10}$–$5\times10^8$ cm$^{-3}$) in helium plasma cooled to cryogenic temperatures ($T_a \lesssim 100$°K). A realistic afterglow model is developed, taking into account electron-temperature relaxation and changes in the density of metastable states. Experiment confirms the existence of the quasistationary decay of electron density, which persists in a broad range of initial dicharge conditions, including the magnetic field. Comparison of experiment with theory has led to the elucidation of the role of hot electrons in cryogenic plasma, to a determination of improved values for the constants of the elementary processes, and, in particular, to a verification of the fact that the temperature dependence of the recombination coefficient of $He_3^+$ ions is of the form $T_e^{-1}$.

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## INTRODUCTION

When the temperature of heavy particles in helium plasma is reduced from $T_a \approx 300$°K down to cryogenic values, the absolute values of plasma parameters can be investigated in a new range, and new quantitative relationships can be established between them. For example, in decaying plasma, the electron temperature $T_e$ falls to 20–100°K, the density of metastable states may reach $M \approx 10^{13}$ cm$^{-3}$, and the ratio of $M$-particle and electron densities rises to $M/n \gtrsim 10^2$ ($n$ is the electron density). There is also a qualitative change in the ion composition of cryogenic plasma. The bulk of the ions consists of the polyatomic ions $He_n^+$, where $n \geq 3$. The recombination coefficient for these ions is $\alpha > 10^{-6}$ cm$^3$·sec$^{-1}$ and exceeds the recombination coefficients of $He^+$ and $He_2^+$, which make up the plasma at room temperature, by a large factor. All this is responsible for many of the features that arise both during the decay of cryogenic plasma and in the cryogenic dc discharge.[1,2]

In the final analysis, the values of the leading plasma parameters and the relationships between them must be traceable to elementary processes (and are ultimately responsible for many aspects of collective phenomena). Sufficient material has now emerged from studies of elementary processes in cryogenic helium plasma to enable us to consider the description of plasma decay as a process of simultaneous variation in its

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leading parameters, above all, $n$, $T_e$, and $M$. It is clear that this can be done only on the basis of a particular model of cryogenic afterglow. The distinguishing feature of a model of this kind is that it must take into account, even in the first approximation, the strong influence of each of the parameters $T_e$, $n$, and $M$ on all others. The enhancement of this interdependence under crygoenic conditions is due to the much higher density of hot electrons as compared with the situation prevailing at $T_a \approx 300°K$. These electrons are continuously created with initial energies in the range $\varepsilon_0 = 14-20$ eV in reactions involving the breakup of $M$-particles. The hot electrons continue to heat the cryogenic gas and the system of cold electrons. The strong source of hot electrons ensures that the cryogenic plasma is highly nonisothermal ($T_e/T_a \approx 10$) even during the late stages of the decay process, right up to a few milliseconds. The prolonged relaxation of $T_e$ in its turn constantly regulates the rate of recombinational perturbation of the electron density. Hot electrons provide a substantial contribution to the cryogenic afterglow and to the balance of the total number of electrons. Processes connected with the behavior of hot electrons are thus found to have a dominant influence. In particular, it is essential to consider what happens to hot electrons reaching the walls of the discharge vessel.

In this paper, we use a realistic model of recombinational decay as a basis for the development of analytic solutions of the balance equations describing variations in $n$, $T_e$, and $M$ in cryogenic helium afterglow. The theoretical results are then compared with experimental data on the kinetics of electron density. The structure of the solutions shows that the time dependence of the electron density $n$ is largely determined by the relaxation of $T_e$ and $M$. Comparison of theoretical and experimental results on electron-density kinetics can thus serve as one of the criteria in judging the validity of possible solutions. Additional information has been obtained from measurements on electron density $n$ in magnetic fields. An external magnetic field modifies the lifetime of hot electrons in plasma, and thus controls the degree of heating of the system of cold electrons. This leads to a change in the rate of recombinational decay of cryogenic plasma in magnetic fields.

## 1. METHOD OF MEASURING ELECTRON DENSITY

The plasma electron density was determined by the microwave resonance method. The procedure used to measure the resonance-frequency shifts is described in detail in a previous paper.[3] Conversion of the frequency shifts to absolute values of electron density was performed by the method described in the same paper. The main changes in the experimental setup involved the design of the microwave resonator and vacuum system for introducing and purifying the working gas.

A silver-coated cylindrical resonator, designed to operate in the $TE_{011}$ mode, was excited at $f = 10$ GHz through a matched coaxial input. The resonator could be tuned by moving an end piston capable of suppressing parasitic $TM$ modes. A superconducting solenoid was wound around the cylindric surface of the resonator and its continuation, and was used to measure the electron density in the magnetic field. The solenoid was 16 cm long and the 4-cm resonator was placed at its center. The solenoid was calibrated with the aid of Hall probes, and the working value of the magnetic field could be varied with a 0.5% accuracy up to 10 000 Oe by varying the current. The electron density was not measured in the region of the cyclotron resonance at $H = 3210$ Oe. The resonance frequency was determined with a precision of $\lesssim 0.5$ MHz for characteristic frequency shifts in the region of recombinational decay between a few units and some tens of megahertz.

The cylindrical discharge chamber was made of quartz glass and had a diffusion length $\Lambda = 0.544$ cm ($L = 3.3$ cm, $R = 1.53$ cm). It was directly joined to the vacuum system used to introduce and purify the working gas. The vacuum system was pumped off and held at a pressure not higher than $10^{-8}$ mm Hg. It was cleaned up by a high-frequency discharge in helium. Gaseous helium was drawn directly from the vapor above the liquid in a dewar, and was passed through a special micron filter placed in liquid helium. The gas was then passed through a system of nitrogen traps filled with a suitable absorbent capable of removing $N_2$, $O_2$, and $H_2$ impurities.

The temperature of the walls of the discharge chamber was monitored with a thermocouple during the measurements. When the discharge vessel was cooled with liquid helium, the temperature of the neutral gas within the volume of the plasma was determined from the period of acoustic oscillations observed on the resonance curve of the resonator. These oscillations were spontaneously excited at the onset of gas breakdown by the hf pulse because of the nonadiabatic heating of the cryogenic gas.[1] The pressure in the system was measured by a special oil manometer. The working pressure $p$ was reduced to 293°K. For $p \approx 1$ mm Hg, the pressure was measured to within 20% and this precision was lowest at working pressures up to a few tens of mm Hg. Figures 1 and 2 show the experimental graphs of electron density as a function of time. It is clear that the nonlinear recombinational decay of electron density continues up to $t \approx 10$ ms both in the presence and absence of the magnetic field, and covers the concentration range $n = 10^{10} - 5 \times 10^8$ cm$^{-3}$.

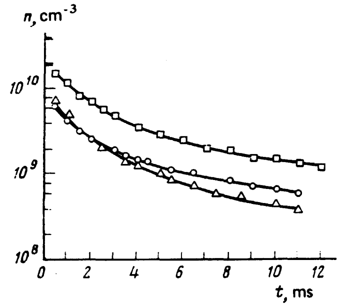


FIG. 1. Decay of electron density ($H = 0$): □—$T_a = 100°K$, $p = 15$ mm Hg; o—$T_a = 8°K$, $p = 17$ mm Hg; △—$T_a = 8°K$, $p = 34$ mm Hg.

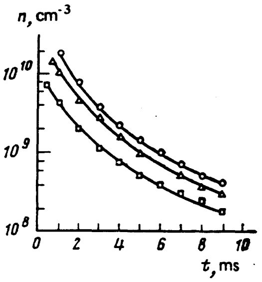


FIG. 2. Decay of electron density for $T_a = 4.2\,°K$, $p = 0.7$ mm Hg; ○—$H = 1320$ Oe, △—$H = 2200$ Oe, □—$H = 0$.

## 2. DECAY MODEL FOR CRYOGENIC PLASMA

It is useful to consider separately a number of groups (systems) of particles in cryogenic helium plasma. Firstly, we take the system of cold electrons of density $n$, which are mainly localized in the energy range between zero and a few $kT_e$ ($T_e = 20–100°K$). We then take the group of hot electrons of density $n^*$ and mean energy $\bar{\varepsilon}^* \gg kT_e$. Quasineutrality is ensured by the ions which, in our experiments (see below), were mainly $He_3^+$. The ion temperature was assumed to be equal to the temperature of the neutral atoms. A further system of particles is made up of the metastables of density $M$. Figure 3 shows the interrelation between the leading elementary processes in which these groups of particles are produced and participate. It also shows schematically the energy flows between these particle systems.

Henceforth, we shall consider the following processes:

$$He_3^+ + e \xrightarrow{\alpha_3} He_2^* + He \begin{cases} \nearrow M_1 + 2He \\ \searrow M_2 + He. \end{cases} \quad (A)$$

$$He(2^3S) + He(2^3S) \begin{cases} \xrightarrow{(1-\chi)\beta} He_2^+ + e^*\,(17.4\text{ eV}), \\ \xrightarrow{\chi\beta} He^+ + He + e^*\,(15\text{ eV}) \end{cases} \quad (B)$$

$$He(2^3S) + e \xrightarrow{\gamma} He + e^*\,(20.6\text{ eV}), \quad (C)$$

$$He_2^+ + 2He \xrightarrow{k_2} He_3^+ + He, \quad (D)$$

$$He_2^+ + e + He \xrightarrow{\alpha_2} He_2^* + He, \quad (E)$$

and the diffusion of $He_3^+$ and $He(2^3S)$. The expressions for the coefficients of these reactions are given in the table. The collision frequencies ($s^{-1}$) are as follows:

$$\delta\nu_{ea} = 3.6\cdot 10^3\, pT_e^{1/2} \text{ for } \varepsilon < 5\text{ eV},\ \delta\nu_{ea} = 7.2\cdot 10^5 p \text{ for } \varepsilon > 5\text{ eV, Ref. 7;}$$

$$\nu_{ee} = 7.7\cdot 10^{-5} n (kT_e\,[\text{eV}])^{-3/2}, \text{ Ref. 7},$$

$$\nu_{ei} = 3.64 n T_e^{-3/2} \ln\Lambda, \quad \Lambda = 1.24\cdot 10^4 (T^3/n)^{1/2}, \text{ Ref. 9}.$$

The metastable states can be characterized as follows. As at $T_a = 300°K$,[10] the equilibrium density $M$ of the triplet metastable states is much higher than the density $S$ of the singlet metastable states. Typically, $M/S > 10$.[1] The rates of reactions of the same type, which involve the participation of these particles, are very similar. We shall therefore confine our attention to the triplet states, i.e., the $M$-particles. The density of these particles consists of the atomic $[He(2^3S)]$ and molecular $[He(2^3\Sigma_u^+)]$ components. The atomic and molecular metastable states can be combined to a single set of levels of the $M$-particles. The energy width of this set (≈2 eV) is much smaller than the level excitation energy (≈20 eV). The rate constants of reactions of the same type involving the participation of metastable atoms and molecules are very similar.[5] The most extensively investigated are reactions involving the participation of atomic states and, therefore, the numerical values of reaction constants involving $M$-particles will henceforth refer to reactions with the participation of $He(2^3S)$.

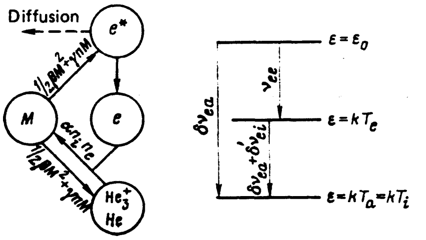


FIG. 3. Schematic representation of the adopted decay model.

Let us consider our decay model in terms of Fig. 3. According to this model, the creation of the $M$-particles occurs in the two-body recombination process between $He_3^+$ and cold electrons [reaction (A)]. The capture of the electron by the ion results in the appearance of the triatomic molecule in a highly excited (autoionization) state. This molecule dissociates along two channels [reaction (A)] and, after cascade transitions, atoms and molecules are produced in the lowest metastable states. The $M$-particles that appear within the working volume decay in the course of mutual collisions [reaction (B)] and during super-elastic collisions with electrons [reaction (C)]. The heavy byproduct of the superelastic decay of the $M$-particles is He in its ground state. The pair reaction (B) can proceed along two channels with the formation of either $He^+$ or $He_2^+$. Even at $T_a = 300°K$, the probability that the $He_2^+$ ion will be created is 0.7. As the relative velocity of the colliding $M$-particles is reduced, this probability rises still further. We shall follow Belevtsev and Mnatsakanyan[11] and assume that the dominant process under cryogenic conditions ($T_a \lesssim 100°K$) is ionization of the $M$-particles with the formation of $He_2^+$ ions.

Recombinational decay will, as before, be determined by $He_3^+$ because $He_2^+$ transforms into $He_3^+$ before it recombines. The conversion of the $He_2^+$ ions occurs with the time constant $(k_2p^2)^{-1}$ in reaction (D). This condition is satisfied subject to the obvious requirement that

TABLE I.

| Process | $T_a$, °K | Values of coefficients |
|---|---|---|
| (A) | 4.2* | $\alpha_3 = \theta/T_e$; $\theta = 4\cdot 10^{-6}$ cm³s⁻¹deg, Ref. 4 |
| | 4.2; 28; 44; 77 | $\theta \sim T_a^{-\xi}$, $\xi = 0{,}32$ Ref. 7 |
| (B) | 300 | $\beta = (4.5\pm 1)\cdot 10^{-9}$ cm³s⁻¹ Ref. 6 $\chi = 0.3$ Ref. 8 |
| | 4.2 ÷ 300 | $\beta \backsim T_a^{1/6}$ Ref. 1 |
| (C) | 300 | $\gamma = 4.5\cdot 10^{-9}$ cm³s⁻¹ Refs. 1, 8 |
| (D) | 80 | $[k_2p^2]^{-1} \ll 9\cdot 10^{-8} p^{-2}$ [s] Ref. 7 |
| (E) | 300 | $\alpha_2 \backsim 1/T_e$; $\alpha_2 = 1{,}1\cdot 10^{-8} + 4.1\cdot 10^{-10} p$ cm³s⁻¹ Ref. 5 |
| Diffusion of $He_3^+$ according to Ref. 3 | ≤100 | $D_a p = 280$ cm²s⁻¹ · mm Hg |
| | 10±3 | $D_a p < 100$ cm²s⁻¹ · mm Hg |
| Diffusion of $He(2^3S)$ according to Ref. 1 | 80 | $D_M p = 200$ cm²s⁻¹ · mm Hg |
| | 5 | $D_M p = 28$ cm²s⁻¹ · mm Hg |

*$n = 5\cdot 10^5 – 2.5\cdot 10^{10}$ cm⁻³.
**$D_M \sim T_a^{2/3}$.

$(k_2p^2)^{-1} \ll 1/\alpha_2 n$, where $\alpha_2$ is the $He_2^+$ recombination coefficient [reaction (*E*)]. If we suppose, for the purposes of rough calculation, that $(k_2p^2)^{-1} = 9\times10^{-3}p^{-2}$ s (which corresponds to the $He^+ \to He_2^+$ conversion at $T_a = 80°K$)[7] we find that this inequality is satisfied for $n \leq 10^{10}$ cm$^{-3}$ and $p\approx 3$ mm Hg. Since the actual value of the time constants for the $He_2^+$ conversion is much smaller (see the table), we may suppose that the conversion process proceeds quite rapidly even for $p\approx 1$ mm Hg.[1)]

According to the experimental results reported by Gerardo and Gusinow,[12] the density ratio at $T_a = 80°K$ is $[He_3^+]/[He_2^+] \gg 1$ for $n = 10^{10}$ cm$^{-3}$ and $p \geq 10$ mm Hg. As the temperature of the gas is reduced, the value of $(k_2p^2)^{-1}$ should fall at least as rapidly as the $He^+$ conversion time constant in the case of the polarization interaction between the atom and the ion, i.e., as $T_a^{3/4}$.[13] Since $[He_3^+]/[He_2^+] \propto k_2p_2$,[12] the $He_3^+$ ions are the main component for $T_a = 4.2°K$ at pressures lower than those for $T_a = 80°K$, i.e., $p\approx 1$ mm Hg.

Thus, the foregoing discussion enables us to select the experimental range of values of $p$ and $n$ in which the pair decay of *M*-particles does, in fact, end in the formation of $He_3^+$ ions, whose equilibrium density provides the dominant component. We note that the working pressures in our experiments were $p \geq 15$ mm Hg for $T_a \leq 100°K$ and $p \geq 1$ mm Hg for $T_a = 4.2°K$.

The above processes of decay of metastable states are accompanied by the appearance not only of the heavy particles, but also hot electrons $e^*$ with initial energies $\varepsilon_0 = 17.4$ and 19.8 eV. These electrons relax from $\varepsilon = \varepsilon_0$ to $\varepsilon \approx kT_e$, thus heating the neutral gas and the system of cold electrons. The characteristic time for cooling these electrons is $\tau = 1/(\delta\nu_{ea} + \nu_{ee})$, where $\delta\nu_{ea}$ is the frequency of the cooling collisions with atoms and $\nu_{ee}$ is the frequency of electron-electron collisions. The relaxation of the hot electrons occurs during their diffusion to the walls of the discharge chamber. The decay of the cryogenic plasma is, in many respects, determined by whether or not the energy level $\varepsilon \approx kT_e$ is populated by hot electrons during the characteristic diffusion time $\tau_d$. If hot electrons reach the surface with $\varepsilon \gg kT_e$, then plasma decay depends on the boundary conditions adopted for them.

For the range of values of $p$, $n$, and $T_a$ used in our experiments, the diffusion of cold electrons, ions, and *M*-particles can be ignored. The condition for the validity of this approximation in the case of electrons and ions is

$$l_R = (D_a/\alpha_3 n)^{1/2} \ll \Lambda_0,$$

where $l_R$ is the mean distance traversed by the charged particle before recombination with $He_3^+$. Using the values of $D_a p$ and $\alpha_3$ given in the table, we find that $l_R \ll \Lambda_0$ even for low pressures $p\approx 1$ mm Hg if $n \gtrsim 5\times10^8$ cm$^{-3}$, $T_e < 50°K$, and $T_a = 4.2°K$. For temperatures $T_a = 80$–$100°K$, we may neglect ambipolar diffusion at pressures exceeding a few mm Hg. These conclusions are confirmed for $n \gtrsim 5\times10^8$ cm$^{-3}$ by the nonlinear form of the measured electron-density curves (see Figs. 1 and 2). By analogy with $l_R$, we can introduce the length $l_M = (D_M/\beta M)^{1/2}$ for the *M*-particles. For $M > 10^{11}$ cm$^{-3}$, which is characteristic for the cryogenic afterglow,[1] diffusion is unimportant under our conditions for $p > 1$ mm Hg, $T_a = 4.2$, and for $p \geq 10$ mm Hg and $T_a \lesssim 100°K$.

It follows from the foregoing that the simplest model can be used to describe plasma decay. In this model, only one type of ion, namely, $He_3^+$, predominates and is responsible for recombination. The diffusion mechanism of decay of cold-particle density is not taken into account. The range of $n$ and $p$ used under these conditions remains sufficiently broad. For example, for $T_a = 4.2°K$, we have $5\times10^8$ cm$^{-3} \lesssim n \lesssim 10^{10}$ cm$^{-3}$ and $p \gtrsim 1$ mm Hg.

Further analysis is complicated by the fact that it is essential to know the fraction of hot electrons reaching the surface of the discharge vessel for a given pressure. It is convenient to consider the following two limiting cases:

$$\tau_d \ll \frac{1}{\delta\nu_{ea}+\nu_{ee}}, \quad \tau_d \gg \frac{1}{\delta\nu_{ea}+\nu_{ee}}.$$

In the first case, which occurs for $p < p_0$, a large fraction of the hot electrons reaches the walls with energies not very different from the initial energy during the characteristic diffusion time $\tau_d$. In the second case, which occurs for $p > p_0$, a large fraction of the hot electrons succeeds in cooling within the working volume and populates the energy level $\varepsilon \approx kT_e$. Since the potential difference between the wall and the center of the discharge chamber is of the order $kT_e/e_0$, hot electrons with $\varepsilon \gg kT_e$ diffuse to the wall under free diffusion conditions with diffusion coefficient $D(\varepsilon) = 2\varepsilon/3m\nu_{ea}$. The pressure $p_0$ can be estimated from the condition $\tau_d(\varepsilon_0) = 1/\delta\nu_{ea}$ ($\nu_{ee} \ll \delta\nu_{ea}$ for $\varepsilon \gg kT_e$). For $\varepsilon_0 = 17.4$ eV, we have

$$\tau_d/\tau_{ea} = 0.92\cdot10^{-1}\Lambda_0^2p^2. \quad (1)$$

In our experiments, $p_0 \approx 6$ mm He. In experiments in which the discharge tube is placed in the waveguide, and must have a smaller radius, the value of $p_0$ will be greater by a substantial factor.

The temperature of the slow electrons is determined by heating through collisions with hot electrons with frequency $\nu_{ee}$, and cooling in collisions with atoms with frequency $\delta\nu_{ea}$. At low pressures, cooling collisions with $He_3^+$ ions with frequency $\delta\nu'_{ei}$ may also provide an appreciable contribution to the energy balance.

The temperature of the heavy plasma particles is assumed equal to the constant quantity $T_a$, determined, for example, from the period of the acoustic oscillations. The approximation $T_a$ = const is justified in the analysis of recombinational decay of electron density by the fact that the coefficient $\alpha_3$ is more dependent on $T_e$ than on $T_a$ (see table). The weak temperature dependence of $\beta$ also enables us to use this approximation. In general, the temperature of the neutral gas is determined by its heating by electrons and cooling by thermal conduction.

## 3. PLASMA DECAY KINETICS FOR $\tau_d/\tau_{ea} \gg 1$

### A. Basic theoretical relationships

The set of equations describing the kinetics of densities $n$, $M$ and energy $kT_e$ can easily be determined with the aid of Fig. 3 when most of the hot electrons succeed in cooling prior to collisions with the chamber walls. The result is

$$\frac{d(n^*+n)}{dt}\approx\frac{dn}{dt}=-\alpha(T_e)n^2+\frac{1}{2}\beta M^2, \quad (2)$$

$$\frac{dM}{dt}=\alpha(T_e)n^2-\beta M^2-\gamma nM, \quad (3)$$

$$\frac{d^3/_2kT_en}{dt}=-\delta\nu_{ea}(T_e)\frac{3kT_en}{2}+\frac{1}{2}\beta M^2\Delta\varepsilon_1+\gamma nM\Delta\varepsilon_2. \quad (4)$$

Here and henceforth, $\alpha\equiv\alpha_3$. The quantity $\Delta\varepsilon$ is the fraction of initial energy expended through the corresponding $M$-particle decay in the heating of the system of cold electrons. It is assumed that electron sources are monoenergetic.

The hot-electron distribution function obtained theoretically for $\tau_d/\tau_{ea}\gg 1$ was used by Blagoev *et al.*[14] to obtain an expression for the effective energy carried off by fast electrons with initial energy $\varepsilon'$ in the cold electron system:

$$\Delta\varepsilon=\int\limits_0^{\varepsilon'}\frac{\nu_{ee}}{\nu_{ee}+\delta\nu_{ea}}d\varepsilon.$$

Since the integrand has a maximum (in the energy range determined by $\nu_{ee}\approx\delta\nu_{ea}$), we can replace the upper limit by infinity and, after integration, obtain

$$\Delta\varepsilon_{1,2}(\text{eV})=2.4\cdot 10^{-5}(n/p)^{1/2}. \quad (5)$$

Apart from numerical factors of the order of unity, this expression is identical with that obtained from the equation $\nu_{ee}(\Delta\varepsilon)=\delta\nu_{ea}(\Delta\varepsilon)$. When the above condition is satisfied for the collision frequencies, the hot electrons begin to experience rapid Maxwellization and the energy $\Delta\varepsilon$ is almost wholly expended in heating.

We must now find the solutions of (2)–(4). At cryogenic temperatures, the characteristic time for a change in the densities $n$, $M$, given by $\tau_n=1/\alpha n$ and $\tau_M=1/\beta M$, respectively, is very different. Since, under cryogenic conditions, $\alpha\gtrsim 10^{-6}$ cm$^3\cdot$s$^{-1}$, $\beta\approx 2\times 10^{-9}$ cm$^3\cdot$s$^{-1}$, $M/n\approx 10^2$, we have $\tau_M/\tau_n\gtrsim 10$. A substantial difference is also found to occur between the characteristic cooling and heating times $1/\delta\nu_{ea}$ and $1/\beta M$ of the system of cold electrons at $T_a=300$°K. For $M\gtrsim 10^{11}$ cm$^{-3}$, $p\gtrsim 1$ mm Hg and $T_e\lesssim 300$°K, we have $\delta\nu_{ea}/\beta M\gtrsim 100$. In principle, this enables us to assume in advance that the approximate solution of (2) and (4) can be obtained by writing these equations in the quasistationary form.[2)]

For a rigorous proof of the existence of quasistationary decay modes, we must identify a small parameter $\mu\ll 1$ in front of the derivatives in (2) and (4). The equation with a small parameter in front of the highest-order derivative belongs to a class with singular perturbation, and its solution has the following properties.[15] In a certain neighborhood of the initial time $t=t_0\equiv 0$, the exact solutions are very different from the degenerate, i.e., those with $\mu=0$ in the so-called boundary layer of width $t_b$. Well away from $t_b$ ($t\gg t_b$), solutions of any accuracy in $\mu$ can be sought in the approximation with $\mu=0$, i.e., in the quasistationary approximation. In particular, if we assume $T_e=$const for (2) and (3), and transform to the dimensionless variables

$$x=\frac{n}{k_1},\quad y=\frac{m}{k_2},\quad \tau=\frac{t}{k_3},\quad k_3=\frac{1}{\beta k_2},$$

we obtain

$$\left(\frac{\beta}{\alpha}\right)^{1/2}\frac{dx}{d\tau}=-x^2+\frac{1}{2}y^2, \quad (6)$$

$$\frac{dy}{d\tau}=x^2-y^2-\frac{\gamma}{(\alpha\beta)^{1/2}}xy. \quad (7)$$

In these expressions, $(\beta/\alpha)^{1/2}=\mu_1\ll 1$ is a small parameter of the problem. We note that, for $T_a=300$°K, the recombination coefficient is $\alpha\approx 10^{-8}-10^{-9}$ cm$^3\cdot$s$^{-1}$ and $\mu_1\to 1$.

According to general theory,[15] the solution of the problem can be sought in the form

$$x(\tau_1\mu_1)=\bar{x}_0(\tau)+\mu_1\bar{x}_1(\tau)+\ldots+\Pi_0x(t^*)+\mu_1\Pi_1x,$$
$$t^*=\tau/\mu_1.$$

There is an analogous solution for $y(\tau_1\mu_1)$. The boundary-layer functions $\Pi x$ and $\Pi y$ need be taken into account only near $t=0$ because they usually decay exponentially.

For $t\gg t_b$, the solutions of (6) and (7) in the first approximation in $\mu_1$ is

$$\bar{x}_0(\tau)=\frac{\bar{y}_0(\tau)}{\sqrt{2}},\quad \bar{y}_0(\tau)=\frac{\bar{y}(0)}{[1+(1/2+\mu_2/\sqrt{2})\bar{y}(0)\tau]}, \quad (8)$$

where $\mu^2=\gamma/(\alpha\beta)^{1/2}$, and $\bar{y}(0)=y(0)$.

Following the algorithm described in Chapter 3 of the book by Vasil'eva and Butuzov,[15] we can set up the differential equation for the function $\Pi_0x$ as well. The expression $t_b=2/(\alpha\beta)^{1/2}M_0$ for the width of the boundary layer is then obtained from the condition for the decay of $\Pi_0x$ by the factor $\exp(2\sqrt{2})$. Since, in cryogenic plasma, $M_0\geqslant 10^{12}$ cm$^{-3}$, we have $t_b<5\times 10^{-5}$ s. The parameter $\mu_2$ enters the original equations in a regular fashion and, therefore, has no effect on the character of the results. Using the values of $\alpha$, $\beta$, and $\gamma$ in the table, we find that $\mu_2\ll 1$. Hence, superelastic decay of $M$-particles need not be taken into account in cryogenic plasmas. At room temperatures, $\mu_2\to 1$.

The equations given by (8), written in terms of the dimensional coordinates $1/n$, $1/M$ as functions of $t$, take the form of straight lines whose slope yields the effective recombination coefficients $\alpha_{\text{eff}}=(\alpha\beta)^{1/2}/2$ and the effective coefficients $\beta_{\text{eff}}=\beta/2$. The recombination coefficient of the ion $He_3^+$ was first measured as $\alpha_{\text{eff}}$ by Fugol' *et al.*[3] It turned out to be $\alpha_{\text{eff}}=1.3\times 10^{-7}$ cm$^3\cdot$s$^{-1}$ at $T_a=10$°K. The result $\beta_{\text{eff}}=1.1\times 10^{-9}$ cm$^3\cdot$s$^{-1}$ at $T_a=10$°K was reported by Fugol' *et al.*[16] The result $\beta=2\beta_{\text{eff}}=2.2\times 10^{-9}$ cm$^3\cdot$s$^{-1}$ is in good agreement with the value $\beta=2.5\times 10^{-9}$ cm$^3\cdot$s$^{-1}$, given in the table at the same temperature. The actual value $\alpha=1.6\times 10^{-5}$ cm$^3\cdot$s$^{-1}$ is very close to the value of $\alpha_3$ in the table if $T_e\lesssim 40$°K. These values of the electron temperature are characteristic (see below) for cryogenic plasmas at the pressures $p \geqslant 10$ mm Hg and densities $n \approx 10^9$ cm$^3$, used by Fugol' *et al.*[3] in their measurements of $\alpha_{eff}$. It follows from the foregoing that the determination of the rate constants for elementary reactions through the analysis of the decay curve for cryogenic plasmas is far from equivalent to measurements of the actual parameters.

Returning now to the original set of equations given by (2)–(4), we note that the quasistationary decay of density $n$ remains valid even for $T_e \neq$ const. The quasistationary variation of the electron energy is established in a time of the order of $1/\delta\nu_{ea}$. In general, the energy-balance equation is also in the class of equations with a small parameter in front of the derivative. Quasistationary solutions of the energy-balance equations are possible and are frequently used in the analysis of plasma decay with $T_a = 300°$K (see, for example, Ref. 1). As a result, the original set of equations for these decay times $t \gg 2/(\alpha\beta)^{1/2}M_0, 1/\delta\nu_{ea}$ includes two algebraic equations and one differential equation. If we solve them and use the data on $\alpha(T_e)$, $\delta\nu_{ea}(T_e)$, given in Sec. 2, together with (5), we obtain

$$M=\frac{M_0}{1+{}^1/{}_2\beta M_0 t},\quad \frac{M}{n}=10.2\left(\frac{\theta p}{n}\right)^{0.3}\beta^{-0.5},$$
$$T_e=1.9\cdot 10^{-2}\theta^{0.4}(n/p)^{0.6}, \qquad (9)$$
$$n^{-0.7}=\text{const}+5.13p^{0.3}\theta^{0.3}\beta^{0.5}t\,.$$

According to these results, for $p \geqslant 10$ mm Hg and $T_a = 4.2°$K, the electron temperature varies from $T_e \leqslant 210°$K for $n = 10^{10}$ cm$^{-3}$ to $T_e \leqslant 52°$K for $n = 10^9$ cm$^{-3}$. The density ratio varies at the same time from $M/n \gtrsim 86$ to $M/n \geqslant 170$. It is easily shown that, under the above conditions, cooling collisions with the ions can be neglected.

### B. Comparison of theoretical results with experimental data

A comparison will not be carried out with the measured electron-density kinetics. The experimental results involved are those obtained for decay times $t \geqslant 1$ msec, i.e., much greater than the characteristic times $t_b$ for the establishment of the quasistationary variation in density $n$. The electron density was measured as a function of time for pressures $p \geqslant 17$ mm Hg and $T_a = 8 \pm 2°$K. At these pressures, $\tau_d/\tau_{ea} \geqslant 8$ according to (1). It may therefore be considered, to a good approximation, that most of the hot electrons remain within the volume and relax to $\varepsilon \approx kT_e$.

The experimental results were then used to determine the function $n(t)$, the slope $\Delta$ of the linear relationship $n^{-0.7} \propto \Delta t$ and the influence of gas pressure and temperature on the slope. Figure 4 shows the measured results in the form of $n^{-0.7}$ as a function of $t$. It is clear that, up to $t = 10$ msec, the decay of electron density is satisfactorily described by the above linear formula [see (9)]. At subsequent times, when $n \lesssim 5 \times 10^8$ cm$^{-3}$, the kinetics begins to be influenced by the diffusion of charged particles, which our model does not take into account. We note that, if we do not take into account the relaxation of $T_e$, then, according to (8), the

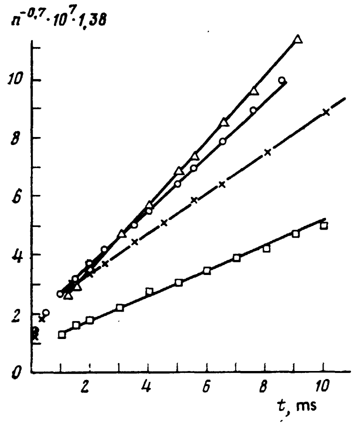


FIG. 4. Electron density ($n$, cm$^{-3}$) as a function of time for $\tau_d/\tau_{ea} \gg 1$. For $T_a = 8°$K: ×—$p = 17$ mm Hg, ○—34.1 mm Hg, △—68 mm Hg, □—$T_a = 100$ °K, $p = 15$ mm Hg.

time dependence is described by $n^{-1} \propto t$. Thus, the relaxation of $T_t$ has a considerable effect on the function $n(t)$.

Let us now compare the experimental and theoretical values of the slope $\Delta$ for $T_a = 8°$K. Using the tabulated values of the coefficients $\theta$ and $\beta$ as functions of temperature, we find that

$$\theta=3.2\cdot 10^{-4}\ \text{cm}^3\text{s}^{-1}\text{deg}\,,\qquad \beta=2.46\cdot 10^{-9}\ \text{cm}^3\text{s}^{-1}\,.$$

According to (9), the calculated slope is $\Delta = 2.27 \times 10^{-5}p^{0.3}\text{s}^{-1}\cdot\text{cm}^{2.1}$. Figure 5 shows the measured slope $\Delta_e$ as a function of $p^{0.3}$. The function $\Delta_e = 2.4 \times 10^{-5}p^{0.3}\text{s}^{-1}\cdot\text{cm}^{2.1}$, shown by the broken line in Fig. 5, lies within the limits of experimental uncertainty. Comparison of $\Delta_e$ and $\Delta$ shows that the theory is in adequate agreement with experiment.

The temperature dependence of the slope of $\Delta \propto T^{-0.013}$ follows from (9) when the temperature dependence of $\theta$ and $\beta$ is taken into account. This very weak dependence ensures that $\Delta$ falls by a factor of only 1.031 between $T_a = 100°$K and $T_a = 8°$K. This cannot be detected experimentally. The function $n(t)$ is satisfactorily represented by $n^{-0.7} \propto t$ even for $T_a \leqslant 100°$K if the discharge is produced by pulses of enhanced hf power.

## 4. PLASMA DECAY KINETICS FOR $\tau_d/\tau_{ea} \ll 1$

In this case, hot electrons do not succeed in cooling down during the diffusion time $\tau_d$, and reach the walls of the discharge chamber with energies not very different from the initial energy. Ions that are produced simultaneously with hot electrons in reaction ($B$) with practically zero energy remain in the discharge volume. The plasma remains quasineutral if one of the following possible conditions is satisfied on the boundary:

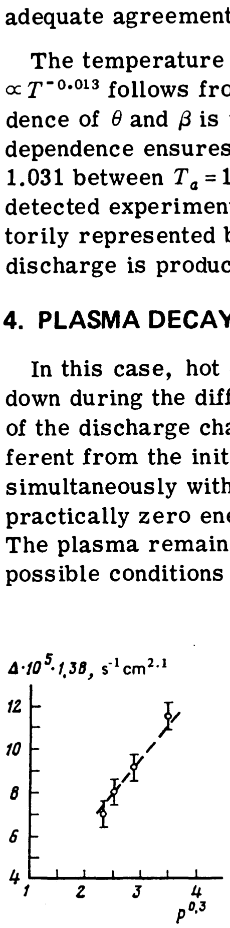


FIG. 5. Measured slope $\Delta$ as a function of pressure ($\tau_d/\tau_{ea} \gg 1$).

I. Hot electrons are reflected elastically ($\varepsilon_{\text{refl}} \gtrsim \Delta\varepsilon_1$) from the surface, i.e., most of them are cooled down within the body of the plasma.

II. Hot electrons are reflected inelastically and, after reflection, have energies between zero and a few $kT_e$.

III. Hot electrons remain on the chamber walls and an ion current flows out of the plasma and compensates the diffusion current of hot electrons.

Let us consider each of these cases separately.

It is clear that, if boundary condition I is realized, the plasma decay process should proceed in the same way as at high pressures, for which $\tau_d/\tau_{ea} \gg 1$ [see (9), Sec. 3]. However, experiments performed for $\tau_d/\tau_{ea} \approx 10^{-2}$ show that (see below) there is an appreciable departure from $n^{-0.7} \propto t$. Moreover, the external magnetic field will produce a change in the rate of recombinational decay of the density $n$, whereas (9) does not predict this. All this enables us to suppose that boundary condition I is not satisfied.

For the next two conditions, the surface of the chamber acts as a sink of hot-electron energy. These electrons heat the system of cold electrons only during the time $\tau_d$. The quasistationary energy-balance equation then assumes the usual form

$$^3/_2 kT_e n(\delta'\nu_{ei}+\delta\nu_{ea}) = \Delta\varepsilon_3 \beta M^2/2. \tag{10}$$

This takes into account cooling electron-ion collisions which are important at low pressures. The quantity $\Delta\varepsilon_3$ is the fraction of initial energy expended in heating. It is given by[14]:

$$\Delta\varepsilon_3 = \tau_d(\varepsilon_1)\nu_{ee}(\varepsilon_1)(\varepsilon_1/\varepsilon_0)^{3/2}\varepsilon_0,$$

where $\varepsilon_1 = 19.8$ eV is the excitation energy of the state $He(2^3S)$. Using the expressions for $\tau_d$ and $\nu_{ee}$, and the fact that $\varepsilon_0 = 17.4$ eV, we find that

$$\Delta\varepsilon_3\,(\text{eV}) = 2.33\cdot 10^{-12} pn\Lambda_0^2. \tag{11}$$

The effect of the magnetic field on the decay process can be taken into account by replacing $\Lambda_0$ with the magnetic diffusion length

$$\frac{1}{\Lambda_H^2} = \left(\frac{2.4}{R}\right)^2 \frac{1}{1+(\omega_H/\nu_{ea})^2} + \left(\frac{\pi}{L}\right)^2, \tag{12}$$

where $\omega_H = 1.76\times 10^7 H$ is the electron-cyclotron frequency.

The quantity $\delta'$ in (10) is the inelasticity parameter for collisions between electrons and $He_3^+$ ions, evaluated per collision with momentum transfer. In plasma, $\delta' < 2m_e/m_i$, since the cross section $Q_1$ for the scattering of an electron by an ion with fractional energy transfer $\delta_e \geqslant 2m_e/m_i$ is less than the total cross section $Q_2$ for scattering with momentum transfer. It is readily shown on the basis of the material given by Mitchner and Kruger[9] that $Q_2/Q_1 \approx \ln\Lambda$. The density balance equations corresponding to boundary conditions II and III can be generalized by introducing the parameter $c$ which is either 0 or 1:

$$dn/dt = -\alpha(T_e)n^2 + c\beta M^2/2, \tag{13}$$

$$dM/dt = \alpha(T_e)n^2 - \beta M^2. \tag{14}$$

When $c = 1$, the hot electrons populate the energy level $\varepsilon \approx kT_e$, which corresponds to boundary condition II. For boundary condition III, $c = 0$.

### A. Basic theoretical relationships in the case of boundary condition II

For $c = 1$, plasma decay is described by (10), (13), and (14). It is easily seen that the approximation of quasistationary variation of the density $n$ remains valid. However, the inclusion of cooling collisions with ions leads to great difficulties in finding analytic solutions of the resulting set of equations. We shall give solutions in the approximation in which $\delta\nu_{ei} \ll \delta\nu_{ea}$. In our experiments, we considered the region of densities $n$ in which plasma decay could be described by these solutions. The influence of electron-ion collisions on plasma decay was examined only in a qualitative fashion. If we solve [taking (11) into account] the initial equation for this case, we obtain:

$$M = \frac{M_0}{1+{}^1/_2\beta M_0 t}, \quad \frac{M}{n} = 1.02\cdot 10^{2.4}\Lambda_0^{-0.4}\beta^{-0.5}n^{-0.4}\theta^{0.3}, \tag{15}$$

$$T_e = 1.9\cdot 10^{-4.8}\theta^{0.4}\Lambda_0 \quad n^{0.8}, \tag{16}$$

$$n^{-0.6} = \text{const} + 0.513\cdot 10^{2.4}\theta^{0.3}\Lambda_0^{-0.4}\beta^{0.5}t. \tag{17}$$

We note at once that, when cooling collisions with ions are taken into account, the exponent $x$ in the relation $n^{-x} \propto t$ is reduced. This can be simply demonstrated by considering the initial equations for the limiting case when $\delta'\nu_{ei} \gg \delta\nu_{ea}$. Thus, as the cooling collisions with ions become increasingly important, the function $n^{-x} \propto t$ with $x < 0.6$ provides a better approximation than (17).

### B. Basic theoretical relationships in the case of boundary condition III

For $c = 0$, plasma decay is described by (10), (13), and (14). Since there is no source of electrons in this variant, and the source of $M$-particles is retained, it can be shown that, for decay times $t \gg 1/\alpha n$, we must have $\beta M^2 \gg \alpha n^2$. This can be shown by combining (13) and (14) into an inhomogeneous differential equation in $M/n$ and then solving it. In this approximation, and for $\delta\nu_{ei} \ll \delta\nu_{ea}$, the solution of the initial equations for the density $n$ is

$$n^{-3/7} \sim 1.37\cdot 10^3\theta^{3/7}\Lambda_0^{-4/7}\beta^{3/7}t. \tag{18}$$

Here, as before, the inclusion of collisions with ions results in a smaller value of the exponent in the expression for the density $n$.

### C. Comparison of theory with experiment

In our experiment, the density $n$ was measured with and without the magnetic field. For the comparison with theory, we used experimental results obtained for decay times $t \geqslant 1$ ms for which the quasistationary state was definitely known to prevail. The measurements were performed in the pressure range $p = 0.7$–$1$ mm Hg. According to (1), under these conditions, $\tau_d/\tau_{ea} \approx 10^{-2}$, i.e., most of the hot electrons reached the walls of the chamber. The chamber was cooled only with liquid helium. The temperature of the gas, determined from

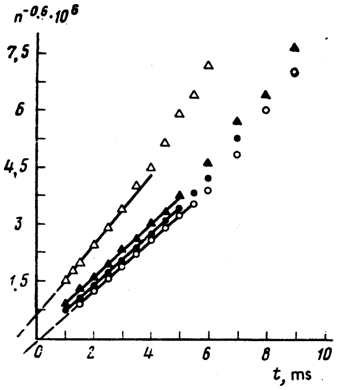


FIG. 6. Electron density ($n$, cm$^{-3}$) as a function of time for boundary condition II, $T_a = 4.2$ °K, $p = 0.7$ mm Hg; △—$H=0$, ▲—$H = 2200$ Oe, ●—$H = 1760$ Oe, ○—$H = 1320$ Oe.

the period of acoustic oscillations, was 4.2°K. The values of the coefficients $\theta$ and $\beta$ for this temperature can be determined from the table. Graphs of $n^{-x} \propto \Delta t$ constructed for different $x$ were then used to determine the slopes $\Delta$ and the dependence of the slope on the magnetic field was established. Figures 6 and 7 plot $n^{-x}$ as a function of $t$, where $x = 0.6$ (boundary condition II), and $x = 3/7$ (boundary condition III).

It is clear from Fig. 6 that, for $n > 10^9$ cm$^{-3}$ and decay times up to a few milliseconds, the experimental results are satisfactorily described by the linear formula $n^{-0.6} \propto \Delta t$. According to (16), the theoretical result in the absence of the magnetic field is $\Delta_0 = 7.3 \times 10^{-4}$ sec$^{-1}$·cm$^{1.3}$, which is in agreement with the experimental result $\Delta_{0e} = (11 \pm 1) \times 10^{-4}$ s$^{-1}$·cm$^{1.3}$. The maximum discrepancy does not exceed 40%. It is readily shown that this discrepancy decreases when the diffusion process is taken into account in the balance equation for the $M$ particles. Diffusion provides a finite contribution when cooling collisions with ions have to be taken into account.

A departure from the above law is observed at the alter stages of the decay process when the concentration is $n < 10^9$ cm$^{-3}$. For $n \approx 2 \times 10^8$ cm$^{-3}$, for which diffusion becomes important, the time dependence can be described by a formula close to $n^{-3/7} \propto t$, which is shown in Fig. 7. According to (18), the theoretical prediction for this dependence with $h = 0$ is $\Delta = 0.3$ s$^{-1}$·cm$^{9/7}$. This is very different (see Fig. 7, broken line) from the experimental results $\Delta_e = (2 \pm 1) \times 10^{-2}$ s$^{-1}$·cm$^{9/7}$. The difference is so great that, to compensate it, one would have to change by two or more orders of magnitude the reaction rates and collision frequencies used in the theory. This could not be justified on any realistic grounds.

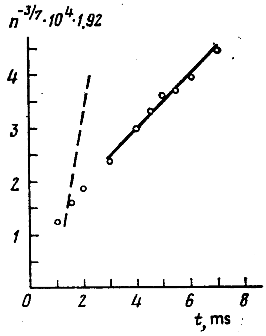


FIG. 7. Electron density ($n$, cm$^{-3}$) for boundary condition III, $T_a = 4.2$ °K, $p = 0.7$ mm Hg, $H = 0$.

The fact that plasma decay is described by the linear formula with $x \approx 3/7$ for $n < 10^9$ cm$^{-3}$ is not unexpected. It is noted above that, when cooling collisions between electrons and ions are taken into account, the exponent $x$ is reduced, and $x < 0.6$ in the case of the boundary condition II. As the electron density decreases, collisions with $He_3^+$ ions are found to compete increasingly with electron-atom collisions. This follows from the result $\delta' \nu_{ei} / \delta \nu_{ea} \sim 1/n^{3/5} \lesssim 1$ ($n \approx 10^{10}$ cm$^{-3}$, $\delta' = \delta$), which is obtained if we use the relationship between $T_e$ and $n$, given by (16). Since the more accurate value is $\delta' \approx \delta / \ln\Lambda$, the dominant effect of cooling collisions with ions on the decay process begins for $n \approx 10^9$ cm$^{-3}$.

It is clear from Fig. 2 that the initial density of electrons in the discharge is higher when the external magnetic field is imposed. The result of this is that the influence of the electron-ion collisions becomes apparent at later stages of the decay process, as compared with $H = 0$. The law $n^{-0.6} \propto t$ is therefore satisfactorily observed (see Fig. 6) in a broader temperature range.

Measurements of $\Delta_H$ performed in magnetic fields confirm the fact that boundary condition II is observed during the decay of cryogenic plasma. The influence of the magnetic field on recombinational decay of plasma can be taken into account by using the magnetic diffusion length $\Lambda_H$ in (17). Since the diameter of the discharge chamber is comparable with its length in our experiments, it follows from (12) that for $H \gtrsim 1300$ Oe, the magnetic diffusion length reaches its limiting value $\Lambda_H = L/\pi$. We then find from (17) that

$$\Delta_0/\Delta_H = (\Lambda_H/\Lambda_0)^{0.4} = 1.3.$$

Figure 8 shows measured values of the slope $\Delta_H$ in different magnetic fields. It is clear that the average experimental value $(\Delta_0/\Delta_H)_e = 1.29$ is in good agreement with the theoretical prediction.

Experiment thus shows that, when $\tau_d/\tau_{ea} \ll 1$, the decay of cryogenic plasma is in good agreement with theoretical predictions if we suppose that condition II is satisfied for the hot electrons. The equations given in Sec. 4A can be used for concentrations $N > 10^9$ cm$^{-3}$ for which electron-ion collisions and diffusion are unimportant. Much more complicated analysis of the initial set of equations (10), (13), and (14) is necessary when the latter effects are taken into account.

## CONCLUSIONS

The agreement between experimental and theoretical results on electron-density kinetics which we have

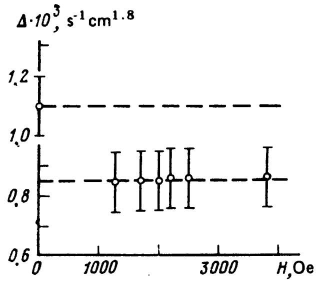


FIG. 8. Measured slope $\Delta$ as a function of magnetic field ($\tau_d/\tau_{ea} \ll 1$).

achieved suggests that the theoretical relationships used in the present work will be suitable for the description of cryogenic afterglow. On the other hand, experiment confirms the validity of the coefficients $\beta$ and $\alpha(T_e)$ used in the theoretical analysis of the decay process.

An important feature of decay in cryogenic plasma is that hot electrons reaching the discharge-chamber walls nevertheless occupy the energy level $\varepsilon \approx kT_e$ within the volume, i.e., they contribute to the overall cold-electron density balance. This ensures that the density balance between $n$ and $M$ remains constant in a broad range of pressures. The relaxation of electron temperature is very dependent on the density ratio of two groups of hot electrons, i.e., electrons reaching the walls and electrons succeeding in relaxing within the volume of the chamber. The internal magnetic field controls this density ratio and, hence, the electron temperature. The magnetic field has its greatest effect when the ratio of the linear dimensions of the discharge chamber is $R/L \ll 1$, and the pressures are chosen so that $\tau_d/\tau_{ea} \ll 1$.

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1)We note that, if we follow Deloche *et al.*[8] and assume that $\alpha_2 \lesssim 10^{-9}$ cm$^3 \cdot$ s$^{-1}$, this inequality is satisfied for much lower pressures.

2)A simple proof consists of solving (2) or (4) on the assumption that the particle and energy sources are strictly constant.

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# Boundary relaxation of a vibrationally excited gas

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The relaxation of a nonequilibrium, vibrationally excited gas, consisting of symmetric diatomic molecules, due to collision with the wall enclosing the gas, is considered. Collisions with the walls remove the vibrational excitation of the molecules. The vibrational energy is distributed among the kinetic degrees of freedom of the molecule and the surface, which can be simulated by a set of harmonic oscillators. An equation is derived for the time variation of the temperature $T$ of the kinetic degrees of freedom of the gas. When $T$ is close to the temperature of the surface $T_s$ and to the population temperature of the vibrational state $T_v$, this equation reduces to the relaxation type. The characteristic time $\tau$ of boundary relaxation depends on the parameters of the gas and of the surface. It is shown that in a certain range of the parameters, boundary relaxation may be more important than the usual relaxation due to pair collisions.

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The total energy of a molecule consists of the energy of the translational, rotational, vibrational and electronic motions. Under conditions of thermal equilibrium, the energy distribution of the molecules with respect to these degrees of freedom is characterized by a single quantity—the temperature. It is not possible to describe a noneqiulibrium molecular gas by a single temperature. The comparatively weak coupling between

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