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VOLUME DISCHARGE IN A GAS, EXCITED BY AN ELECTRON BEAM UNDER CONDITIONS OF NONUNIFORM IONIZATION

O. B. Evdokimov, G. A. Mesiats, and V. B. Ponomarev

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16. Abstract

Volume discharge in gases excited by electron beams is discussed. The steady state electric field distribution is derived. The voltage-current and energy characteristics are described.

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VOLUME DISCHARGE IN A GAS, EXCITED BY AN ELECTRON BEAM UNDER CONDITIONS OF NONUNIFORM IONIZATION /357*

O. B. Evdokimov, G. A. Mesiats, and V. B. Ponomarev**

I. INTRODUCTION

Volume discharge in gases excited by electron beams [1] has found wide applications in powerful lasers [2] in powerful commutators [3], plasmatrons, etc. The theory of such discharges assumes that the distribution in the rate of generation of the electron-ion pairs is uniform along the depth of the discharge gap. However, in reality this is not always the case. Because of the probabilistic nature of electron scattering in gases, the spread in the energy spectrum of the electrons upon their passage through the foil and into the gas, and also because of time-variations in the voltage of the accelerating diode, the distribution of the generation rate of the electron-ion pairs can turn out to be non-uniform. In effect, the beam electrons become thermalized [4]. This leads to nonuniform concentration of plasma in the discharge column and to a nonuniform electric field, which in certain cases leads to instabilities in volume discharge [5]. Since these instabilities are connected with the special properties of the injected beam, they are called injection instabilities [5]. In the case of lasers, the nonuniform ionization leads, in addition, to a nonuniform population inversion which is manifested in the optical steady-state properties of the laser. The question of distribution of the electric field under

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the conditions of nonuniform ionization is discussed in part and on a qualitative level in [6]. In contrast with [6] we will derive the steady-state electric field distribution taking into account shock ionization. We will also consider the voltage-current and energy characteristics and will take up the question of how the discharge properties, at short-time intervals, of injected beams are affected by the nonuniformity in the rate of degeneration of the electron-ion pairs.

We will start with the following system of equations for the one-dimensional case:

$$\frac{\partial n_-}{\partial t}-\frac{\partial n_- v_-}{\partial x}=\Psi+\alpha n_- v_- -\beta n_- n_+ + g,$$
$$\frac{\partial n_+}{\partial t}+\frac{\partial n_+ v_+}{\partial x}=\Psi+\alpha n_- v_- -\beta n_- n_+, \qquad (1)$$
$$\frac{\partial E}{\partial x}=\frac{e}{\varepsilon}(n_+ - n_-),\quad v_\pm=\mu_\pm E.$$

Here, Ψ is the volume rate of generation of electron-ion pairs, the number of electron-ion pairs created by the fast electron beam by ionization of atoms per second per unit volume; g is the volume speed of injection of beam electrons as a result of thermalization; $n_\pm$, $u_\pm$, $v_\pm$ are the concentration, mobility, and speed of the ions and electrons respectively; e is the electron charge; E the electric field intensity; β - recombination coefficient; α - the coefficient of shock ionization. Here and in what follows, the anode is located at x = 0 and the cathode - at x = d. /358

By order of magnitude $g/\Psi \sim \varepsilon_I/\langle T\rangle$, where ε_I is the average energy of electron-ion pairs created, $\langle T\rangle$ is the average energy of fast electrons. We assume that the thermalization time of fast electrons is short.

II. Electric Field Distribution in the Steady State

One of the remarkable properties of volume discharge at large Ψ is the small value of the cathode voltage drop U_k, in comparison with the applied voltage U_o, $U_k << U_o$ [7]. Let us consider a discharge column where we neglect ion drift. In this case, our system of equations (1) becomes:

$$\Psi(x)+g(x)+\alpha\mu_-n_-E=\beta n_+n_-, \tag{2}$$

$$j=j_0(x)+e\mu_-n_-E, \quad \frac{dj_0}{dx}=eg(x), \tag{3}$$

$$\frac{\partial E}{\partial x}=\frac{e(n_+-n_-)}{\varepsilon}, \tag{4}$$

where j_o is the beam current density through the gas at a depth x. In the steady-state case, usually $j_o(x) << e\mu_-n_-E$; hence, $j \sim e\mu_-n_-E$. Let us suppose, just as in the case of uniform ionization, that in the discharge column n_+ is approximately equal to n_-, i.e.,

$$|n_+-n_-|\ll n_-, \tag{5}$$

Then, by eliminating the electron concentration from (3), we obtain in an implicity form the field distribution

$$e\Psi(x)+\alpha(E)j=\frac{\beta j^2}{e\mu^2E^2}, \tag{6}$$

where we have set $\mu_- \equiv \mu$. Figure 1 gives examples of typical distributions $\Psi(x)$ obtained by solving the kinetic equation which describes the transport of monoenergetic electrons in nitrogen in a uniform field and by taking into consideration the separating foil. The fast electrons are injected through the cathode. The function $\Psi(x)$ decreases from the cathode to the anode; nowhere is it equal to zero, since an electron beam propagating through the medium is accompanied by brehmsstrahlung. Approximately $\Psi_{min}/\Psi_{max} \sim 10^{-2} \div 10^{-3}$. With increasing field intensity $\Psi(x)$ becomes more uniform.

Clearly, starting with some minimum value $E(x = xm) = E_{min}$ near the cathode, the field increases in the direction of the anode. If $\frac{d\Psi}{dx}>0$ for all x, then x_m is at the boundary between the discharge column and the cathode region. When the discharge voltage is several times smaller than the static breakdown voltage, then at the point where the electric field is a minimum, the inequality $\Psi(x_m) >> \alpha(E_{min})j/e$ holds and therefore

$$E^2_{min}\approx\frac{\beta J^2}{e^2\mu^2\Psi(x_m)} \tag{7}$$

and by combining (6) and (7) for $0 < x < x_m$, we have:

$$\frac{\Psi(x)}{\Psi(x_m)} = \frac{E_{min}^2}{E^2(x)} - \frac{\mu E_{min}\alpha(E)}{\sqrt{\beta\Psi(x_m)}}. \quad (8)$$

It is possible to show, by considering (8) and (4), that condition (5) is true if /359

$$\frac{\varepsilon E_{min}\sqrt{\beta\Psi(x_m)}}{2e\Psi^2(x)}\left|\frac{d\Psi}{dx}\right| \ll 1. \quad (9)$$

Condition (9) takes into account the fact that $E<E_{min}\sqrt{\Psi(x_m)/\Psi(x)}$. Hence, whenever shock ionization is significant, condition (9) is an overestimate and is fulfilled for realistic values of $\Psi(x)$ (see Figure 1). Figure 2 shows an example of a calculation of the field E(x) for discharge in nitrogen. It shows the influence of shock ionization on the distribution of the field. The curve labeled 2 is a calculation without shock ionization. Curve 3 is a calculation taking shock ionization into account according to equation (8), with parameters μ = 500 cm^2/sec; $\beta = 2 \times 10^{-7}$, the $\alpha(E)$ dependence is taken from [8].

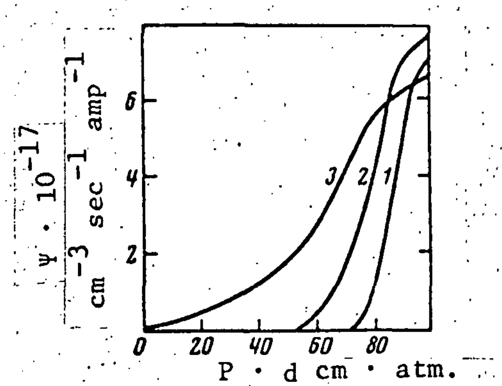

Figure 1: The distribution of the generation speed of electron-ion pairs in nitrogen at atmospheric pressure behind an aluminum foil of 50 μ thickness. Curve 1 - input energy 150 keV, E/p $\approx$ 4 kV/atm·cm; Curve 2 - energy 200 keV E/p = 2 kV/atm·cm; Curve 3 - energy 200 keV, E/p = 4 kV/atm·cm at the beam point $5 \cdot 10^{-4}$ a/cm^2.

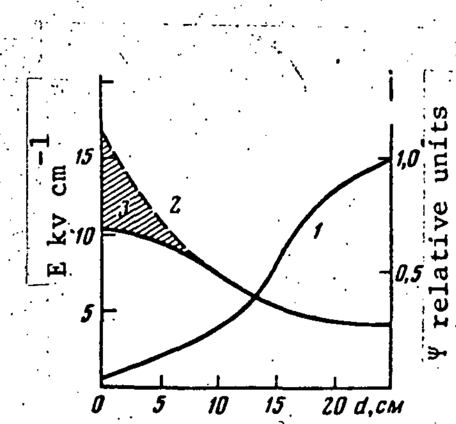

Figure 2: The distribution of the electron-ion pair generation rate (curve 1) in relative units in nitrogen at atmospheric pressure at beam current $5 \cdot 10^{-4}$a, initial energy 200 keV, foil thickness 50 μ; 2- electric field distribution without taking into account shock ionization; 3- electric field distribution with shock ionization.

The calculation uses the distribution $\Psi(x)$ (curve 1) calculated for a uniform field $E = E_{min}$. Since $E(x) \geq E_{min}$, the real distribution is a smoother function and hence, because of Equation (8), curves 2 and 3 give an upper limit for the field distribution.

3. CURRENT-VOLTAGE AND ENERGY CHARACTERISTICS

The relationships (6), (7), and (8) allow us to calculate the current voltage (CVC) and the energy characteristic of volume gas discharge with inhomogeneous ionization with the condition $\int^{x_m} E(x)dx \approx U_0$, where U_0 is the applied gap voltage. If we neglect shock ionization, then from (6) it follows that

$$E(x) = \frac{j\sqrt{\beta}}{e\mu\sqrt{\Psi(x)}}. \qquad (10)$$

Integrating (10) from $x = 0$ to x_m, we obtain CVC /360

$$j = \frac{e\mu U_0}{\sqrt{\beta} x_m \langle \Psi^{-1/2} \rangle}; \qquad (11)$$

here,

$$\langle \Psi^{-1/2} \rangle = \frac{1}{x_m} \int_0^{x_m} \frac{dx}{\sqrt{\Psi(x)}}. \qquad (12)$$

When U_0/d is increased, the distribution $\Psi(x)$ far away from the cathode becomes more uniform (Figure 1) and according to (11) and (12) for sufficiently large distances between the electrodes, we can expect that with increasing U_0, the current will increase more rapidly than linearly.

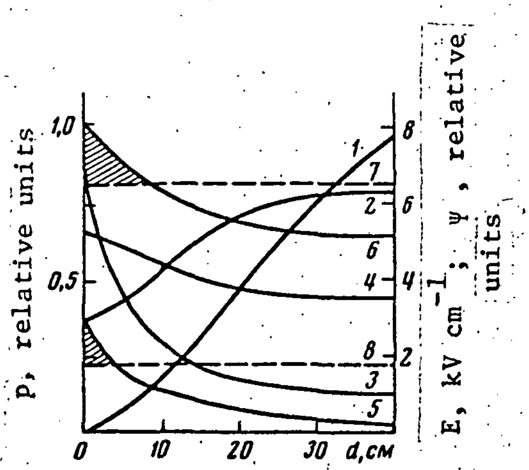

Figure 3: Distribution of generation rate $\Psi(x)$ (curves 1 and 2) electric field (curves 3 and 4) specific power (curves 5 and 6) in nitrogen, at atmospheric pressure, behind a foil of 50 μ thickness for electric field intensities of 2 and 4 kv/cm, respectively; initial beam energy 200 keV, beam current $5 \cdot 10^4$ amp/cm^2. Dashed lines (7, 8) denote the levels of specific power corresponding to the case of uniform ionization.

Specific power generated in the gas discharge upon neglecting shock ionization is calculated from:

$$P_H = jE = \frac{e\mu U_0^2}{d^2\sqrt{\beta\Psi(x)}}\left\langle\frac{1}{\sqrt{\Psi}}\right\rangle^{-2}, \quad x_m \approx d. \tag{13}$$

Equation (13) shows that specific power generated in the discharge increases with decreasing Ψ; i.e., in the direction of the anode. In the case of uniform ionization $\Psi(x) = \Psi_0$, specific power generated in the discharge is equal to

$$P_0 = \frac{e\mu U_0^2\sqrt{\Psi_0}}{d^2\sqrt{\beta}} \tag{14}$$

From (13) and (14) we obtain,

$$\frac{P_H}{P_0} = \frac{1}{\sqrt{\Psi_0\Psi(x)}}\left\langle\frac{1}{\sqrt{\Psi}}\right\rangle^{-2}. \tag{15}$$

If one assumes that $\Psi_0 \sim \Psi\ (x = d)$, then $\sqrt{\Psi_0} > \left\langle\frac{1}{\sqrt{\Psi}}\right\rangle^{-1}$; hence, we can expect in the vicinity of the anode that $P_H > P_0$. In Figure 3 at current density $5 \cdot 10^4$ a/cm^2 initial electron energy 200 keV, aluminum foil thickness 50 μm, U = 80 and 160 kv respectively, we show the distribution Ψ(x) curves 1, 2, E(x) - curves 3, 4, P_H - curves 5, 6. Dashed lines 7 and 8 show the level of specific power P_0. For given U_0 and d, shock ionization is negligibly small. From Figure 3, it is apparent that in the vicinity of the anode we not only get an increase in the field, but also in the generated specific power. Moreover, $P_H > P_0$, in spite of the fact that /361

$$\frac{\int_0^{x_m} P_H(x)\,dx}{\int_0^{x_m} P_0\,dx} = \frac{1}{\sqrt{\Psi_0}}\left\langle\frac{1}{\sqrt{\Psi}}\right\rangle^{-1} < 1. \tag{16}$$

4. SMALL TIME NON-STATIONARY SOLUTIONS

For short times $t < \tau_p = \frac{1}{2\sqrt{\Psi\beta}}$, and large Ψ we can neglect the drift and the recombination terms in (1) and in the discharge column we have $n_- \approx n_+ \approx \Psi t$.

The drift term is negligible if $\left|\frac{\partial n_- v_-}{\partial x}\right| \ll \Psi(x)$ for $t = \tau_p$, which is equivalent to the inequality

$$\frac{\mu E_{min}\sqrt{\Psi_{min}}}{4\sqrt{\beta}\Psi^2}\left|\frac{d\Psi}{dx}\right| \ll 1, \tag{17}$$

where we set $t = \tau_p$. The system of equations (1) with the condition (17) gives an equation for the charge ρ:

$$\frac{\partial\rho}{\partial t}+\frac{e\mu\Psi t}{\varepsilon}\rho=eg+e\mu Et\frac{\partial\Psi}{\partial x}. \tag{18}$$

From the right-hand side of (18), we see that the volume charge is created at the expense of the thermalized electrons g and the gradient of the generation rate of the current carriers. In the regions where we can neglect the creation of volume charge at the expense of $d\Psi/dx$, we obtain from (18):

$$\rho=egt\int_0^1\exp[-at^2(1-u)]du, \quad a=\frac{e\mu\Psi}{2\varepsilon}. \tag{19}$$

The time dependence of the volume charge density ρ shows an extremum. For $at^2 \ll 1$, the charge density grows in time as

$$\rho\approx egt\left(1-\frac{2}{3}at^2\right),$$

and for large t (but smaller τ_p), we obtain

$$\rho\sim\frac{eg}{2at},$$

where ρ is a decreasing function of time.

Let us estimate the maximum ρ_m. The function

$$f=t\int_0^1\exp[-at^2(1-u^2)]du$$

is limited by the inequality

$$t\exp(-at^2)<f<[1-\exp(-at^2)]a^{-1}t^{-1},\ t>0.$$

Hence, the maximum value

$$\max(te^{-at^2})<\max f<\max\left(\frac{1-\exp(-at^2)}{at}\right).$$

The left-hand side of the inequality has the extremum at $at^2 = 0.5$ and the right-hand side at $at^2 \sim 1.25$. Hence, $0.43\ a^{-1/2} < \max f <$

$0.64 \cdot a^{-1/2}$ is the extremum value of the density $\rho_m \cong \frac{eg}{2\sqrt{a}}$ and the evaluation of the maximum field at the expense of thermalized electrons is

$$E_T \approx \frac{1}{2}\int_0^d \frac{\rho_m(x)}{\varepsilon}dx = \frac{e}{4}\int_0^d \frac{g(x)\,dx}{\sqrt{a(x)}} \sim \frac{j_0 d}{4R\langle\sqrt{a}\rangle\varepsilon}, \quad R>d. \tag{20}$$

For beam currents $j_0 = 1\ a/cm^2$, $\varepsilon_I = 35$ ev, $\beta = 2 \cdot 10^{-7}\ cm^3/sec$, we have: $\tau_p = 5 \cdot 10^{-8}$ sec, $\sqrt{a} = 5 \cdot 10^8\ sec^{-1}$, experimental time $a^{-1/2} = 2 \cdot 10^{-9}$ sec, and for distances of the order of $R \sim d$, we obtain $E_T \sim 5$ kv/cm. Because $a \sim j_0$, it follows from (20) that $E_T \sim \sqrt{j_0}$. The above estimate shows that, under the conditions considered, the contribution of thermalized electrons to the field is appreciable for currents of the order of 1 amp and higher. Moreover, in a time proportional to $a^{-1/2}$ the volume charge density undergoes relaxation asymptotically as t^{-1}. The ratio of the maximum value of charge density of thermalized electrons ρ_m to their saturation density ρ_H is independent of g and is determined by the species of gas under consideration

$$\frac{\rho_m}{\rho_H} = \sqrt{\frac{e\mu}{\varepsilon\beta}}.$$

The relaxation is connected with the time growth of conductivity (if $a^{-1/2} \ll \tau_p$).

Let us consider next the building up of volume charge neglecting thermalized electrons. Let us use the equation for the total current which follows from (1)

$$\frac{\partial E}{\partial t} + \frac{e\mu}{\varepsilon} n(x,t)E(x,t) = \frac{1}{\varepsilon} J(t). \tag{21}$$

Assuming that the voltage along the discharge column remains practically constant, after integration with respect to x, equation (21) gives

$$J(t) = \frac{e\mu}{x_m}\int_0^{x_m} n(x,t)E(x,t)\,dx. \tag{22}$$

Given that at high electron pair generation rates

$$n(x,t) = \sqrt{\frac{\Psi(x)}{\beta}}\,\mathrm{th}\, t\sqrt{\Psi\beta} \tag{23}$$

and assuming that $\Psi(x)$ is determined only by the average field U_0/d, and therefore is time-independent, from (21), (22), and (23), we obtain an integral equation for the field:

$$E(x,t)=\frac{U}{d}e^{-\gamma}+\frac{e\mu}{\varepsilon\sqrt{\beta}\,x_m}e^{-\gamma}\int_0^{x_m}\sqrt{\Psi'}\,dx'\int_0^t e^{\gamma'}E(x',t')\,\mathrm{th}\,t\sqrt{\Psi'\beta}\,dt', \tag{24}$$

where $\gamma'=\gamma(t',x')=\nu\ln\mathrm{ch}\,t'\sqrt{\Psi'\beta}$, $\nu=e\mu/\varepsilon\beta$, $\Psi'=\Psi(x')$.

Equation (24) can be solved by iteration. As the zero-order approximation, let us take $E = U_0/d$. Then for $t \to \infty$, it follows from (24) that

$$\frac{E_1(x,t)}{E}\cong\frac{\int_0^{x_m}\sqrt{\Psi}\,\mathrm{th}\,t\sqrt{\Psi\beta}\,dx}{x_m\sqrt{\Psi(x)}}\Bigg|_{t\to\infty}=\frac{\langle\sqrt{\Psi}\rangle}{\sqrt{\Psi(x)}}. \tag{25}$$

As in equation (8), here $E(x, \infty) \sim 1/\sqrt{\Psi(x)}$. The results of (25) agree with (8) to within the constant coefficient, and are valid if $\langle\sqrt{\Psi}\rangle\approx\left\langle\frac{1}{\sqrt{\Psi}}\right\rangle^{-1}$. For small times when $n(x, t) = \Psi(x)t$ from (24), we obtain

$$\frac{E_1(x,t)}{E}=\frac{\langle\Psi\rangle}{\Psi(x)}+\left(1-\frac{\langle\Psi\rangle}{\Psi(x)}\right)\exp(-at^2), \tag{26}$$

i.e., for $at^2 \gg 1$, but $t < \tau_p$:

$$\frac{E_1(x,t)}{E}\approx\frac{\langle\Psi\rangle}{\Psi(x)}. \tag{27}$$

Comparison of (27) and (25) shows that the electric field intensity in the discharge column goes through an extremum.

As is apparent from (25), the relaxation of the field from its extremum to the stationary value is determined by recombination and takes place in a time $t \sim (2 - 3)\,\tau_p$. The field distribution at the time of the extremum can be calculated exactly. For $at^2 \gg 1$, $\tau_p > t$, $n(x, t) = \Psi(x) \cdot t$, we can take $\frac{\partial E}{\partial t}\approx 0$ and then it follows from (21) that

$$E^{\text{эк}}=\frac{c}{\Psi(x)},$$

where c is a constant determined by the condition $\int_0^{x_m} E\,dx\approx U_0$.

Therefore,

$$E^{\text{эк}} = \frac{E}{\Psi(x)}\left\langle\frac{1}{\Psi}\right\rangle^{-1}.$$

This result is more exact, the better inequality $\tau_p \gg a^{-1/2}$ is satisfied.

Specific power at extremum field (27) is determined by:

$$P_{\text{н}}^{\text{эк}} = \frac{e\mu \bar{E}^2 t_{\text{эк}}}{\Psi(x)}\left\langle\frac{1}{\Psi}\right\rangle^{-2} \approx \frac{\sqrt{2e\mu\varepsilon}\,\bar{E}^2}{\Psi^{1/2}(x)}\left\langle\frac{1}{\Psi}\right\rangle^{-2},$$

when $at^2 \gg 1$, since

$$j = e\mu n E = e\mu\Psi t_{\text{эк}} E^{\text{эк}} = e\mu t_{\text{эк}} E\left\langle\frac{1}{\Psi}\right\rangle^{-1}. \tag{28}$$

Thus, short-time considerations show that for non-uniform ionization after a time of the order of $\sim a^{-1/2}$, the field and the specific power undergo a significant redistribution within the gas volume. Afterwards, during the time (2 - 3) τ_p, one sees a decrease in inhomogeneity. Since the first stage takes place in a relatively short time, afterwards, conditions are especially favorable for those mechanisms of instability of volume gas discharge which are determined by the field intensity rather than specific power.

We must remark that (28) does not exclude the possibility that $P_{\text{н}}^{\text{эк}} > P_{\text{н}}\ (t \to \infty)$.

5. EXPERIMENTAL RESULTS

The fact that during electron injection the electric field in the gas gap can be strongly distorted is experimentally verified by the observation of instability of the gas discharge when the energy spectrum of the injected electrons is broadened [9, 10]. For injection pulse times of 10^{-4} sec., the increase in the fall-off of the voltage at the top of the accelerating voltage pulse always led with other variables held constant to the instability of the discharge [9]. The instability first appears in the anode region, i.e., specifically at the location where field increases must occur as a result of redistribution of the potential caused by the inhomogeneous distribution of conductivity in the gap. We must remark that the large duration of the pulse front of the accelerating voltage can likewise be a cause of instability.

The effect of the spectral characteristic of the electron beam of duration 10^{-8} sec. on the instability growth time was investigated in [10]. It was shown that lowering the maximum energy of the injected electrons from 180 keV (current, 50 amps) to a value for which the current was 1 amp, sharply reduced the instability growth time for discharge in nitrogen at 7 atmospheres. In a number of cases this reduction was by a factor of 100 and meant a $t_n \sim 10^{-9}$ sec. Such short instability growth times are caused by streamer discharge initiated in the region of field increase, for which [11]

$$pt_n \approx \frac{\ln N_{\text{кр}}}{\alpha v_-}, \tag{29}$$

where p is the gas pressure and $N_{\text{кр}} \approx 10^8$ is the critical electron number in the avalanche. At nitrogen pressures of p = 7 atm and $t_n \sim 10^{-9}$ sec., the electrical field in the region of inhomogeneity growth must substantially exceed 10^5 volts/cm. Such a field can be easily due to thermalized, low-energy electrons [5]. We should note that formula (29) may be useful in the investigation of fields due to electrons thermalized in gases.

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