A SIMPLE THEORY OF BREAKDOWN OF MONATOMIC NONLIGHT GASES IN FIELDS OF ANY FREQUENCY FROM LOW TO OPTICAL

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Simple, compact, and universal equations are derived for the estimate of the frequencies of ionization and threshold breakdown fields in monatomic not-too-light gases. The equations have a large range of applicability; they are valid for fields of any frequency: constant field, high-frequency, microwave, optical (i.e., for the description of laser breakdown), and for any pressure, duration of the action of the field, and dimension. The equation for the breakdown threshold ensures reasonable agreement with experiment and can be useful for estimates in those cases where experimental data are not available.

A number of works [1-13] (also see the review on laser flash in [14, 15]) have been devoted to the avalanche theory of breakdown of gases by optical radiation, its formulation, solution of the equations, elucidation of the physical details, refinement of the mathematical statements, and computations of the threshold fields. In [1] the theory was based on the quantum kinetic equation for the energy distribution function of electrons, and it was shown that if the energy of the photon is small compared to the characteristic energy of the electrons, the quantum equation reduces to the classical equation with adequate accuracy, which, in particular, describes the breakdown by microwave radiation. The theory of microwave breakdown has been discussed in detail in the book [16], and in a number of cases it gives good agreement with experiment. However, the computations are generally very complex and are not suitable for a clear physical interpretation of the effects.

In the present article a simple approximate theory of avalanche breakdown is proposed. The obtained equations clearly demonstrate the basic physical characteristics of the development of electron avalanche in electromagnetic fields and the regularities of the breakdown and make it possible to obtain rapid estimates.

The theory put forward here is based on the concepts and approximations of [1], a more correct refined variant of the theory of [1], whose results were given in a short communication without derivation [17] for the explanation of the experimental results on the breakdown of gases by CO_2 laser radiation, presented there. However, compared to the computation in [17] the computation here contains significant improvement.

1. Formulation of the Problem and Simplifications. Following the arguments of [1] about the possibility of description of optical breakdown on the basis of the classical equation, we shall start from the equation for the energy distribution function of electrons $n(\varepsilon, t)$, which we write in the form [1]

$$\frac{\partial n}{\partial t} = -\frac{\partial j}{\partial \varepsilon} + Q, \quad j = -D\frac{\partial n}{\partial \varepsilon} + nu - nu_s \tag{1.1}$$

$$D = 2u\varepsilon = A\varepsilon, \quad A = \frac{2}{3}\frac{e^2E^2}{m}\frac{\nu_m}{\omega^2 + \nu_m^2}, \quad u_s = 2\frac{m}{M}\varepsilon\nu_m \ . \tag{1.2}$$

Here j is the "flux" of electrons along the energy axis ε, D is the corresponding "diffusion" coefficient, E is the root-mean-square electric field of the electromagnetic wave $E = E_0/\sqrt{2}$, where E_0 is the amplitude, ω is the frequency of the field, ν_m is the effective frequency of elastic collisions of electrons with atoms, u_s is the "velocity" determined by elastic losses, m and M are masses of electrons and atoms,

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and Q includes terms describing inelastic collisions, production of electrons, and losses connected with the diffusion drift of electrons from the region of action of the field.

The distribution function is normalized to the density of the electrons:

$$\int_0^\infty n(\varepsilon, t)\, d\varepsilon = N_e(t) . \tag{1.3}$$

In the general form the problem is very complex; therefore, we shall restrict the range of investigated conditions and introduce certain simplifications.

1. We shall consider only monatomic gases in order to disregard the excitation of molecular oscillations and the excitation of lower electronic levels.

2. We shall consider only not-too-light gases in order to neglect elastic energy losses of electrons.

3. The excitation and ionization cross sections of atoms $\sigma^*(\varepsilon)$, $\sigma_i(\varepsilon)$ grow from zero at the corresponding potentials I^*, I on increasing the electron energy ε. We introduce the quantities I_1^*, I_1, which slightly exceed (approximately by 1 eV) I^*, I. We neglect the corresponding collisions for $\varepsilon < I_1^*$, I_1, and the frequency of exciting collisions in the interval $I_1^* < \varepsilon < I_1$; $\nu^* = N_a v^* \sigma^*$ is assumed constant (N_a is the density of atoms, v^*, σ^* are some mean velocity of electron and excitation cross section for the given interval).

4. We assume that the electrons, having attained energy I_1, instantaneously experience inelastic collision, ionizing an atom with probability β and exciting it with a probability $1 - \beta$.

5. We assume that electrons undergoing inelastic collisions and also electrons knocked out from atoms as a result of ionization appear with "zero" energy, even though actually they have a small energy.

6. We assume the frequency of elastic collisions ν_m and the characteristic time of diffusion drift of electrons from the region of action of the field τ_d to be constant. The frequency of the diffusion drifts is $\nu_d = \tau_d^{-1} = D/\Lambda^2$, where D is the coefficient of free diffusion of electrons in ordinary space and Λ is the diffusion length of the order of the characteristic dimension of the region.†

With these approximations Eq. (1.1) becomes

$$\frac{\partial n}{\partial t} = -\frac{\partial j}{\partial \varepsilon} - \delta \nu^* n - \nu_d n, \qquad j = -D\frac{\partial n}{\partial \varepsilon} + nu , \tag{1.4}$$

$$\delta(\varepsilon) = \begin{cases} 0 & \text{for } \varepsilon < I_1^* \\ 1 & \text{for } I_1^* < \varepsilon < I_1 . \end{cases} \tag{1.5}$$

The action of the inelastic collisions for $\varepsilon > I_1$ and the sources of electrons in the range of small energies are replaced by the corresponding boundary conditions. For $\varepsilon = I_1$ there is an infinite power "sink" and, hence,

$$n = 0 \text{ for } \varepsilon = I_1 . \tag{1.6}$$

By definition $j(\varepsilon, t)$ is the number of electrons per cm^3, which cross the point ε on the energy axis per second; $j_1 = j(I_1, t)$ electrons reach the "sink" at $\varepsilon = I_1$ per second and slowly lose their energy up to "zero," while a fraction $j_1\beta$ of these gets doubled. Besides, electrons which experienced inelastic collisions in the energy range $I_1^* < \varepsilon < I_1$ are produced with zero energy. Thus the rate of production or the flux $j_0 = j(0, t)$ is

$$j_0 = j_1(1 - \beta) + 2j_1\beta + \nu^* \int_{I_1^*}^{I_1} n\, d\varepsilon . \tag{1.7}$$

This is also the second boundary condition.

At the boundary of the region when $\varepsilon = I_1^*$, where the function $\delta(\varepsilon)$ is discontinuous according to equation (1.5), n and j are continuous, i.e., $\partial n/\partial \varepsilon$ is continuous.

† For a cylinder of diameter d and height L, $1/\Lambda^2 = (\pi/d)^2 + (2.4/L)^2$; for a sphere, $\Lambda = d/2\pi$.

In this approximation the average ionization energy over the entire spectrum

$$\nu_i = \frac{1}{N_e}\int_0^\infty n v \sigma_i(\varepsilon)\, d\varepsilon \qquad (1.8)$$

is determined by the equality $\nu_i N_e = j_1\beta$. If Eq. (1.4) is integrated over the entire spectrum taking account of this last equality and Eqs. (1.6) and (1.7), we obtain the phenomenological kinetic equation

$$dN_e/dt = \nu_i N_e - \nu_d N_e = N_e / \Theta, \quad 1/\Theta = \nu_i - \nu_d \ . \qquad (1.9)$$

In a field of constant amplitude the stationary spectrum is established very soon, and it is clear from (1.9) that the solution of Eq. (1.4) should be sought in the form

$$n(\varepsilon, t) = n(\varepsilon)\exp(t/\Theta)$$

where Θ is the as yet unknown time constant of avalanche, which is related to the unknown ionization frequency ν_i through (1.9).

The function $n(\varepsilon)$ satisfies the equation

$$(\nu_i + \delta\nu^*) n + \frac{dj}{d\varepsilon} = 0, \quad j = -D\frac{dn}{d\varepsilon} + nu \qquad (1.10)$$

and the boundary conditions (1.6) and (1.7). Since the equation is homogeneous, one of the boundary conditions is "superfluous" and this offers the possibility of determining the functions $\nu_i(E)$ and $\Theta(E)$.

We shall call the appearance of N_1 electrons in the gas after a time t_1 of the action of the electromagnetic pulse for N_0 initial (seed) electrons the "breakdown." Then obviously the threshold field E is determined from the condition

$$1/\Theta(E) = \nu_i(E) - \nu_d = t_1^{-1}\ln N_1/N_0 \ . \qquad (1.11)$$

In the case of sufficiently long pulses ($t_1 \to \infty$), we obtain the "stationary" criterion of the breakdown:

$$\nu_i = \nu_d \qquad (\Theta = \infty) \ .$$

2. Solution of the Equations and Ionization Frequency. The system of equations (1.10), (1.6), (1.7) can be solved exactly. The general solution of (1.10) has the form

$$n = C_1 \exp\left[2\sqrt{(\nu_i + \delta\nu^*)\varepsilon/A}\right] + C_2\exp\left[-2\sqrt{(\nu_i + \delta\nu^*)\varepsilon/A}\right] \ . \qquad (2.1)$$

Subjecting it to the boundary conditions, we obtain a system of four linear equations for the four constants of integration $C_{1,2}$ (in the two regions). Equating the determinant of this system to zero, we obtain the transcendental equation for the unknown ν_i.

Omitting the simple but quite long intermediate operations, we give this equation in the final dimensionless form

$$e^{(a-1)y}\left(\operatorname{ch}\frac{y}{z} + z\operatorname{sh}\frac{y}{z}\right) - e^{-(a-1)y}\left(\operatorname{ch}\frac{y}{z} - z\operatorname{sh}\frac{y}{z}\right) =$$
$$= 2a(1+\beta)y + 2(1 - z^{-2})\{y\operatorname{ch}[(a-1)y] + \operatorname{sh}[(a-1)y] - ay\} \qquad (2.2)$$

where we have introduced the following notation:

$$y = 2\sqrt{(\nu_i + \nu^*) I_1^*/A} = a^{-1}\sqrt{6(\nu_i + \nu^*)/\nu_E}$$
$$z = \sqrt{(\nu_i + \nu^*)/\nu_i} = \sqrt{1 + \nu^*/\nu_i}, \qquad a = \sqrt{I_1/I_1^*}, \qquad \nu_E = 3A/2I_1 \ . \qquad (2.3)$$

The quantity $\nu_E = 1/\tau_E$ represents the "frequency of the energy sets" because τ_E is the time required, according to the elementary theory, for the electrons in the field to acquire energy I_1 sufficient for multiplication. We recall that according to the elementary theory $d\varepsilon/dt = 3A/2$.

Numerically,

$$\nu_E = \frac{1.75\cdot 10^{15}E^2}{\omega^2 + \nu_m^2}\frac{\nu_m}{I_1} = \frac{6.34\cdot 10^{17}S}{\omega^2 + \nu_m^2}\frac{\nu_m}{I_1} \qquad (2.4)$$

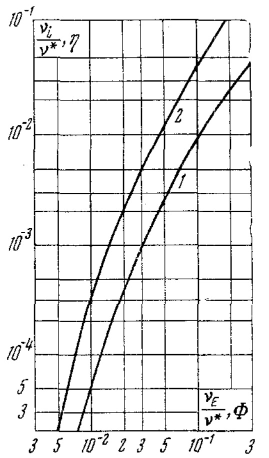

Fig. 1

where $S = cE^2/4\pi$ is the energy flux density in the electromagnetic wave, E in V/cm, I_1 in eV, ω, ν_m in sec^{-1}, and S in W/cm^2.

Equation (2.2) together with equations (2.3) determines the dimensionless dependence of ν_i/ν^* on ν_E/ν^*, which contains two parameters a and β. The parameter a is practically the same for most monatomic gases; it can be taken equal to 1.2. An investigation of the excitation and ionization cross sections of a majority of atoms shows that β can also be taken to be equal for all atoms and we can put $\beta = 0.2$. However, one significant fact should be noted here. Electrons with energies between the excitation and ionization potentials mainly excite the lower excited levels, and this results in energy losses of the electrons (in any case, at the frequency of neodymium laser and lower frequencies). Electrons with energies $\varepsilon > I_1$ excite mainly the upper states. During the breakdown of gases by ruby and neodymium laser radiations ($\hbar\omega = 1.78$ and 1.17 eV) these states are rapidly ionized as a result of one- or two-quanta photoelectric effect, so that the excitation of the upper states, just as the ionization by electron impact, leads to a fast breeding. Therefore, in these cases we can take $\beta = 1$.

Thus Eqs. (2.2), (2.3) give two universal dimensionless functions $\nu_i/\nu^* = F(\nu_E/\nu^*)$ for $\beta = 0.2$ (breakdown in constant field at high and microwave frequencies by longwave infrared radiation of CO_2 laser, $\lambda = 10.6\ \mu$, $\hbar\omega = 0.124$ eV) and $\beta = 1$ (light frequencies). These functions, i.e., the dimensionless ionization frequencies, are plotted in Fig. 1 (ordinate) as functions of the dimensionless frequency of the energy sets, which at the same time serves as the dimensionless function determining the threshold field for the breakdown. Curve 1 is for $\beta = 0.2$, curve 2 for $\beta = 1$.

In the two limiting cases the solutions have a simple form accessible to a clear physical interpretation. If the excitation losses are small, $\nu^* \ll \nu_E$, $z \to 1$, Eq. (2.2) with $\beta = 1$ reduces to the equality $\sinh ay = 2ay$, from which we get $ay = 2.18$ and $\nu_i = 0.8\nu_E$. The ionization frequency practically coincided with the frequency of energy sets, which is quite natural. In particular, this case corresponds to the regime in which all excited atoms are rapidly ionized under the action of the radiation (here I should be replaced by I*). This perhaps occurs in the breakdown of gases by the radiation from a ruby laser (see [1, 7, 12, 14]).

If the excitation losses are significant

$$\nu^* \gg \nu_E, \quad y \gg 1, \quad \nu_i \ll \nu_E, \quad z \gg 1, \quad y/z \approx a^{-1}\sqrt{6\nu_i/\nu_E} \ll 1$$

Eq. (1.3) has the asymptotic solution

$$z^2 \approx y^2 \exp\left[(a-1)\,y\right]/12\,a\beta$$

from which we have

$$\nu_i = \nu_E \alpha\beta a^2, \qquad \alpha = 2a\exp\left(-\frac{a-1}{a}\sqrt{6\nu^*/\nu_E}\right). \tag{2.5}$$

It is easy to verify that the quantity α in this case represents the ratio of the fluxes $j(I_1)/j(I_1^*)$, i.e., nothing else than the probability that an electron "skips" the "unsafe" zone $I_1^* < \varepsilon < I_1$ and attains the energy I_1, when it has a chance of producing multiplication.

Thus, as was to be expected, with an accuracy up to a coefficient $a^2 \sim 1$ the ionization frequency is determined by the frequency of the energy sets and by the probabilities α, β that the electron attains the energy I_1 and after this produces the ionization. The electron executes $k = 1/\alpha\beta$ cycles of motion along the energy axis, picking up and losing energy before it multiplies; accordingly the time required for new generation increases by a factor k.

3. Threshold fields. We denote by Φ the function inverse to F determining the dimensionless ionization frequency. Then we have

$$\nu_E/\nu^* = \Phi(\nu_i/\nu^*).$$

Obviously function $\Phi(\eta)$ gives the same graph as in Fig. 1 if the argument is assumed to be numbers plotted along the ordinate in Fig. 1 and the function is taken as numbers plotted along the abscissa. We write the general condition of breakdown (1.11) in the form

$$\nu_i = \nu_d + \nu_t \quad (\nu_t = t_1^{-1} \ln N_1 / N_0) \; ,$$

then the threshold field is determined from the equality

$$\nu_E / \nu^* = \Phi \left[(\nu_d + \nu_t) / \nu^* \right]$$

which after substituting the expression for ν_E gives

$$E^2 = \frac{I_1 m (\omega^2 + \nu_m^2)}{e^2} \frac{\nu^*}{\nu_m} \Phi(\eta), \qquad \eta = \frac{\nu_d + \nu_t}{\nu^*} \; . \tag{3.1}$$

Estimates show that as the average energy $\bar{\varepsilon}$ in the expression for the spatial diffusion coefficient averaged over the spectrum

$$D = \langle v^2 / 3\, \nu_m(v) \rangle = 2\bar{\varepsilon}/3m\nu_m ,$$

one can take one half of the excitation potential of the atoms: $\bar{\varepsilon} \approx I^*/2$.

Thus, for the computation of the threshold field for the breakdown of different gases at different frequencies, pressures, and dimensions we have a very compact numerical equation:

$$E^2 = 5.7 \cdot 10^{-16} I_1 (\omega^2 + \nu_m^2) \frac{\nu^*}{\nu_m} \Phi(\eta) \tag{3.2}$$
$$\eta = (\nu_d + \nu_t) / \nu^*, \qquad \nu_d = 5.8 \cdot 10^{14} I^* / \nu_m \Lambda^2 \; .$$

Here and below E is in V/cm; I_1 in eV; ω, ν in $\sec^{-1}$; and Λ in cm; the function $\Phi(\eta)$ is given by Fig.1.

The case of "stationary" breakdown, when $\nu_t \ll \nu_d$ and $\nu_i \approx \nu_d$, has a large range of applications. Almost all the cases encountered in practice, i.e., breakdown by constant, high-frequency, microwave fields, and infrared radiation of CO_2 laser, come under this category. In this case the limiting laws of breakdown are valid. It is also possible to determine asymptotic analytical expressions for E. At low pressures p or large frequencies $\omega^2 \gg \nu_m^2$ and when diffusion losses predominate

$$\alpha \approx 1, \; \eta = \nu_d / \nu^* \sim 1 / p^2 \to \infty, \; \Phi \approx \eta / \beta$$

from (3.2), we obtain

$$E \approx \sqrt{\frac{I_1 I^*}{3\beta}} \; \frac{\omega}{\nu_m \Lambda} \sim \frac{\omega}{p\Lambda} \; , \tag{3.3}$$

i.e., E/ω is a function of $p\Lambda$.

At high pressures or low frequencies (in particular in constant field) for Φ we obtain the asymptotic equation

$$\Phi = 6 \left(\frac{a-1}{a} \right)^2 [\ln (2a^3 \beta \Phi / \eta)]^{-2} \, ,$$

and in this case ($p \to \infty$, $\eta \to 0$) we obtain

$$E = \frac{5.8 \cdot 10^{-8} \sqrt{I_1 \nu_m \nu^*} \; (1 - \sqrt{I_1^* / I_1})}{\ln (2a^3 \beta \Phi / \eta)} \tag{3.4}$$

where Φ under the logarithmic symbol can be taken as constant. It is evident then that "breakdown voltage" $E\Lambda$ is a function of $p\Lambda$ (reminiscent of the well-known Paschen curves) and that the asymptotic dependence of the threshold on the pressure for $p \to \infty$ has the form

$$E \sim p \, / \, [\mathrm{const} + \ln (p\Lambda)] \; ,$$

i.e., grows slightly more slowly than p and depends on the dimensions only logarithmically.

4. Comparison with Experiment. By way of illustration we use equation (3.2) for computing the thresholds for breakdown of argon and xenon in different frequency ranges.† For argon we take $\nu_m = 7 \cdot 10^9 p$, $\nu^* = 2.6 \cdot 10^8 p$ (according to the data presented in [16]), $I_1 = 16.8$ eV, $I^* = 11.5$ eV; for xenon, $\nu_m = 1.5 \cdot 10^{10} p$, $\nu^* = 4 \cdot 10^8 p$ (this last value is chosen on the basis of the data in [18]), $I_1 = 13.1$ eV, $I^* = 8.4$ eV; here p is in torr.

†Elastic losses in argon are small, and in xenon they are generally insignificant.

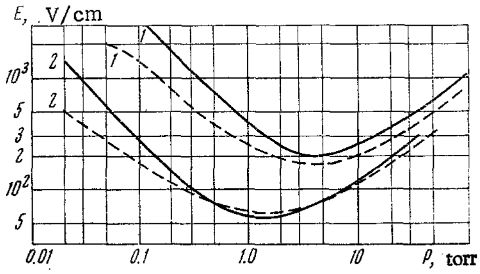

Fig. 2

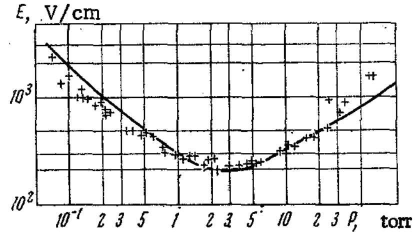

Fig. 3

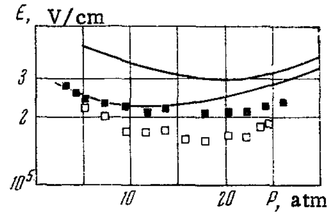

Fig. 4

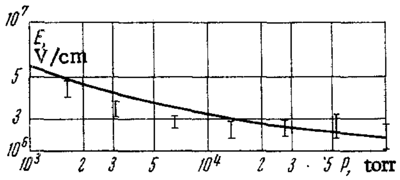

Fig. 5

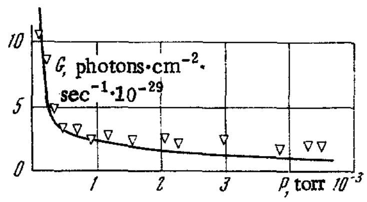

Fig. 6

The threshold fields for microwave breakdown of argon are shown in Fig. 2. The continuous curves pertain to computations, the dashes to experiment. The experimental data are taken from [16]. Curve 1 is for the frequency of the field equal to 2.8 MHz ($\omega = 1.8 \cdot 10^{10}$ rad/sec), $\Lambda = 0.15$ cm; curve 2 for the frequency equal to 0.99 GHz ($\omega = 6.2 \cdot 10^{9}$ rad/sec), $\Lambda = 0.63$ mm. There are no computations for argon in [16]. The threshold fields for microwave breakdown of xenon are shown in Fig. 3. The experimental data were taken from [16]. The frequency of the field is 2.8 MHz, $\Lambda = 0.10$ cm. The continuous curve pertains to the computation. In [16] there are no computations for xenon.

The threshold fields for the breakdown of argon and xenon by the radiation from CO_2 laser ($\lambda = 10.6\,\mu$, $\omega = 1.99 \cdot 10^{14}$ rad/sec) are given in Fig. 4. The experimental data were taken from [17]; the pulse length $t_1 \sim 1\,\mu$sec, the radius of the focus $4 \cdot 10^{-3}$ cm; black squares are for argon, white for xenon. The upper curve pertains to computation for argon, the lower to the computation for xenon. The threshold fields for the breakdown of argon by neodymium laser are shown in Fig. 5. The experimental points are taken from [19]. The pulse length $t_1 = 50$ nsec, $\Lambda = 1.64 \cdot 10^{-3}$ cm. The continuous curve gives the result of the computation. The threshold flux density of photons for the breakdown of xenon by neodymium laser are given in Fig. 6. Here the experimental points are taken from [20]. The pulse length is 35 nsec, and the radius of the focus is $4.5 \cdot 10^{-3}$ cm. The continuous curves pertain to the computation.

Furthermore, on the basis of the general equation for the ionization frequency obtained above, we computed the first ionization coefficient of Townsend for a constant field $\alpha_i = \nu_i/\mu E$, where $\mu = e/m\nu_m$ is the mobility. From (2.5) the asymptotic equation for α_i (at large pressures and small E/p) is

$$\alpha_i = A_i E e^{-Bp/E}, \quad A_i = 2a^3\beta/I_1$$
$$B = \frac{a-1}{a}\sqrt{\frac{6I_1 m\nu_m\nu^*}{e^2p^2}} = 5.8\cdot10^{-8}\frac{a-1}{a}\sqrt{I_1\frac{\nu_m}{p}\frac{\nu^*}{p}} \tag{3.5}$$

where A_i is in ion pairs/V and B is in V/cm · torr. For argon with the same constants B = 53, $A_i = 0.04$. If the experimental curve (graph 4.49 in [21]) is approximated by function (3.5), we obtain B = 31, $A_i = 0.01$. For xenon, the agreement is still better: B = 85, $A_i = 0.05$ from the computation; B = 85, $A_i = 0.1$ from the experiment.

Thus the approximate theory discussed above, which led to a very simple and universal equation (3.2) for the breakdown threshold gives reasonable agreement with the experiment in all frequency ranges from constant field to optical and in a wide range of pressures clearly demonstrating the characteristic regularities of the phenomenon.

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