Dynamic methods of confinement and stabilization of hot plasma

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A review is given of the current state of studies of the confinement and stabilization of hot plasma using high-frequency magnetic fields in the megahertz region. The principles upon which such studies are based are summarized. A theoretical analysis is reported of a plasma confinement system using "traveling-wave" high-frequency fields and of the results of experiments performed with systems of this kind. Another section is devoted to dynamic plasma stabilization systems, the fundamentals of the theory of the dynamic stabilization of the Z pinch, and the results of experiments performed in this field. Experimental data are compared with the basic assumptions of the theory and certain general conclusions are reported with regard to further developments and possibilities in this subject.

CONTENTS

1. INTRODUCTION

Investigations into the properties of hot plasma, which have now been continuing for over twenty years, have, in essence, been concerned with searches for methods of confinement and stabilization of unstable plasma configurations.

The present review is devoted to one of these methods, in which the required effect is achieved by exposing the plasma to high-frequency electromagnetic fields. Although the basic aim of all these studies has been conditioned by the problem of controlled thermonuclear fusion, many of the problems in this area are of more general interest and may turn out to be very significant in other branches of physics.

High-frequency fields can be used to produce field configurations whose average variations per period are such that the equilibrium and stability conditions can be satisfied for plasmas placed in such fields. The period of the field oscillations must be made small enough to ensure that the variation in the field per period does not lead to a substantial deformation of the plasma, or to its intensive interaction with the walls of the gas-discharge chamber. The minimum frequency of the applied field is given by

$$f_{\min} \approx \frac{v_{Ti}}{L}, \qquad (1.1)$$

where v_{Ti} is the thermal velocity of the plasma ions and L is a characteristic linear dimension of the system.

The above formula gives an order-of-magnitude estimate. More accurate expressions for the minimum frequency will be given below for specific systems involving dynamic interaction with the plasma.

We must now define the frequency region which is of interest in connection with the problem of controlled thermonuclear fusion. The thermal velocity of deuterium ions at a temperature of ~10 keV is $v_{Ti} = 10^8$ cm/sec. This is, in fact, the optimum temperature when the DT reaction is employed. If we suppose that L is of the order of 10 cm, we find that $f_{\min} \approx 10^7$ sec^{-1}. The wavelength Λ of the electromagnetic wave which corresponds to this frequency in vacuum is 30 m, which is much greater than the geometric dimensions of the system, so that the electromagnetic field can be looked upon as quasistationary, and the electric field is less than the magnetic field in the ratio of $L/\Lambda \ll 1$. The effect of the electric field can therefore be neglected.

Methods which employ the application of strong electromagnetic fields in this frequency band to plasmas, with a view to confinement and stabilization, will be referred to as dynamic methods. The phrase "strong fields" means that the forces associated with the magnetic pressure due to the action of these fields are comparable with the gas pressure in the hot plasma. The expression "dynamic methods" is introduced because the field period in the above frequency band is comparable with the intrinsic times which are characteristic for macroscopic processes in plasma, namely, the time taken by particles moving with thermal velocities to traverse a given distance, the period connected with the propagation of the ions found in plasma, the period of oscillations at the Alfven frequency, and the growth times of MHD instabilities. The plasma can therefore "feel" the field oscillations, and its configuration tends to adjust to the changes in the configuration of the high-frequency fields. The duration of the entire plasma confinement and stabilization process in the case of dynamic methods is much greater than the period of the applied high-frequency fields.

The present review is restricted to methods involving the application of quasistationary fields to plasmas. Questions connected with the use of wave fields in devices in which the wavelengths are comparable with the dimensions of the system must be the subject of a separate review.

The first ideas on the confinement of plasma by quasistationary high-frequency fields and on the dynamic stabilization of plasma in connection with the problem of controlled thermonuclear fusion were put forward in the late fifties. Experimental studies began in the early sixties. In 1963, Osovets[1] published a review of plasma confinement and stabilization by high-frequency electromagnetic fields. In the present paper, which appears a decade later, we attempt to return to this theme, using more recent ideas.

Existing methods of hot-plasma confinement and stabilization by electromagnetic fields have been the subject of extensive investigations, are well known, and there is little point in describing them in detail here. These are the toroidal traps (Tokamak and Stellarator), open adiabatic traps, and pulsed systems based on the Z and θ pinches. Recent developments involve two new methods based on the utilization of strong laser and electron beams. Each has its own specific merits and difficulties.

We believe that dynamic methods which involve the interaction between the plasma and high-frequency electromagnetic fields will occupy an important place among other methods in connection with the problem of controlled thermonuclear fusion.

High-frequency fields can be used to produce fields with configurations approaching the ideal trap for collisionless plasma. It is then possible to satisfy the conditions of equilibrium and macroscopic stability for values of the parameter $\beta = 8\pi p/H^2$ approaching unity.

This opens up a relatively broad range of different possibilities. We shall not describe here all the existing variations on this theme, and will concentrate our attention on common physical properties of such systems, using specific devices as characteristic examples rather than as final recipes. In the words of the great Czech chess-player, Richard Réti, "variations die but ideas remain."

The basic ideas can be briefly formulated as follows:

a) Plasma confinement—the gas kinetic pressure of the hot plasma is opposed by the pressure associated with the high-frequency magnetic field. It turns out that the field configuration can be made such that the field on the surface of the plasma always increases toward the periphery. The rate of increase of the field can be made as high as desired and is restricted only by energy considerations.

It is, of course, assumed that the frequency of the field variations is sufficiently high [see (1.1)] to ensure that the time-averaged magnetic pressure acts on the plasma. It is probably basically impossible to produce a configuration using only time-independent fields in which the mean H^2 is a maximum everywhere on the surface of the plasma and, at the same time, the field curvature between the boundary of the plasma and the periphery is sufficiently large.

In the case of high-frequency fields, it is possible to ensure that the plasma is everywhere separated from the walls of the discharge chamber. This can be achieved even when the plasma is in a closed volume, for example, in a toroidal chamber.

Since the characteristic length over which the plasma pressure changes from zero to the maximum value is determined by the depth of penetration of the confining field into the plasma, the pressure gradient is observed only over the depth of the skin layer (the confining high-frequency field penetrates as a result of the skin effect to a relatively small depth). This is why the radial plasma pressure distribution is not as smooth as in the case of slowly varying confining fields and is characterized by a steep profile which clearly differentiates the region occupied by the plasma from the surrounding regions. Since a strong high-frequency field is present in the region between the plasma and the wall, this region cannot contain an appreciable concentration of the neutral gas or ionized particles because the neutrals are ionized by strong electric fields whilst charged particles drift into the region occupied by plasma. Thermal insulation of the plasma from the chamber walls in such systems should therefore be better than in systems with constant or relatively slowly varying confining magnetic fields.

These advantages of plasma confinement by high-frequency fields are, however, negated by practical difficulties connected with the technology of devices having parameters approaching those necessary for producing the thermonuclear fusion reaction with a positive energy yield. These difficulties are to do with energy, and will be considered in the next section when we analyze the operation of a specific system of this kind.

b) Plasma stabilization. In this case, the plasma pressure is balanced by the pressure associated with the slowly varying field. The function of the high-frequency field is to maintain stability, which is mainly of the MHD type. This can be achieved by high-frequency fields alone, or in combination with constant or slowly varying fields (these are the hybrid systems).

Here, because of the time variations in the stabilizing-field configuration, the instability does not succeed in developing in any particular direction provided, of course, the stabilizing-field configuration varies at a sufficiently high rate. In effect, there is an alternation of conditions favoring and impeding the development of instabilities, but the plasma deformation is never allowed to reach danger level, and so the plasma remains stable on the average.

Systems of this kind are frequently referred to as systems with dynamic stabilization.

We now begin our review of specific systems by considering high-frequency traps. As an example, we start the "traveling wave" trap.

2. "TRAVELING WAVE", THEORETICAL FOUNDATIONS

In this trap, the field configuration can be such that the mean value of the square of the magnetic field always increases from the center to the periphery. This field can be produced by a set of conductors in which the phase of the currents in neighboring conductors is shifted by an angle θ_T. The simplest example of this system is the time-independent trap with cusp geometry in which this angle is equal to π. In this case, the field at the center of the vessel is zero and increases rapidly toward the periphery. However, in this trap, there are gaps between the conductors through which the particles can escape and reach the chamber walls.

Such traps have been extensively investigated and have now been almost totally rejected because the escape of the particles through the gaps is unacceptably high.

In 1957 a system was proposed [2] in which the gaps were continuously displaced, so that the particles did not succeed in leaving the useful volume. This can be achieved only with the aid of high-frequency fields which are produced by a set of conductors carrying high-frequency currents whose phase is shifted between adjacent conductors by a constant angle θ, greater than zero and less than π. In this system, the amplitude of the magnetic field propagates in the direction of the z axis (Fig. 1). This is the so-called "traveling-wave system." When it is closed into a torus, the wave rotates in a way similar to, for example, the configuration in an electric induction motor.

The wave velocity is $v_{ph} = \omega/\kappa$, where the wave number $\kappa = 2\pi/\lambda$ and λ is the space wavelength. The conditions ensuring that the perturbation of the plasma surface due to the moving gap is small can be reduced to (1.1). If we take the characteristic dimension of the system to be the minor radius of the torus $a = 10-100$ cm, and take the thermal valocity as $v_{Ti} = 10^7-10^8$ cm/sec, then $\omega = 10^7-10^8$ sec^{-1}, which corresponds to a wavelength of 300–30 m, i.e., the wavelength is always much greater than all the characteristic dimensions. We can therefore neglect displacement currents in comparison with conduction currents, and suppose that the system is quasistationary.

This frequency criterion will, of course, ensure that particles with velocities much greater than the thermal velocity, i.e., those from the tail of the plasma velocity distribution function, can in principle succeed in entering the gap. However, the fact that the field increases toward the periphery should prevent appreciable escape of such particles. Moreover, it will be shown below that the quasiconstant magnetic field due to the drift current, which appears in the traveling-wave system, and will also impede the escape of fast ions from the plasma. However, these problems require separate analysis which will not be given here. We merely report the basic relationships characteristic for devices of this kind.

To avoid cluttering the present review with long calculations, we confine our attention to the simplest problems with a view to establishing the physics of the situation. We begin by considering the solution for the plane problem. We have a set of conductors carrying alternating currents in such a way that currents in neighboring conductors are shifted in phase by an angle θ. Strictly speaking, the problem should be considered for a torus, and the corresponding set of equations should be written in terms of toroidal coordinates. The latter problem has been solved (see, for example, [3,4]) but the solution is quite complicated and will not be reproduced here. The solution for the plane case is useful when the distance between the supply coil and the surface of the plasma is much less than the radius of the plasma formation. Since this requirement has been satisfied in the experiments performed so far, we consider this approach to be justified. We shall formulate the problem as follows. Consider the surface of plasma interacting with a set of conductors whose disposition and numeration are shown in Fig. 1. The shortest distance between the conductors and the plasma surface will be denoted by h and the distance between neighboring conductors by b. The plasma is assumed to be perfectly conducting and its surface undisturbed. The conditions which ensure that the perturbation of the surface is sufficiently small will be given below. Under these assumptions, a magnetic field pressure can be assumed to be applied to a plane surface. We now calculate the magnetic field produced by the currents flowing in an infinite number of conductors, arranged as shown in the figure, at a point z on the plasma surface.

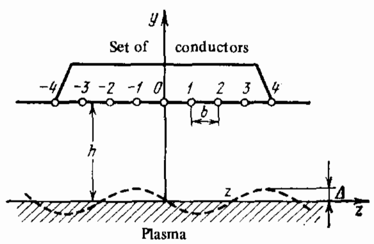


FIG. 1. "Traveling wave" calculations for the plane case in terms of the coordinates y, z. The shaded region represents the plasma. The broken line shows the surface perturbation, where Δ is the perturbation amplitude. The conductors are numbered relative to an arbitrary point; b is the distance between conductors in the supplying coil, h is the distance between the system of conductors and the unperturbed plasma surface.

Using the method of images, we obtain the field at the point z on the surface of the plasma in the form of the series

$$H(z, t) = \mathrm{Re}\frac{4\pi J_0}{bc} e^{i\omega t} \sum_{n=-\infty}^{n=+\infty} \frac{(h/b)\, e^{in\theta}}{(h/b)^2 + [n-(z/b)]^2}; \qquad (2.1)$$

where J_0 is the amplitude of the current in the conductors, θ is the phase difference between currents in neighboring conductors, and n is the number of the conductor.

Evaluating the sum in (2.1), and taking the real part, we find that

$$H(z, t) = \frac{4\pi J_0}{bc} \frac{\mathrm{sh}\,[(h/b)(2\pi-\theta)]\cos[\omega t + (\theta z/b)] + \mathrm{sh}\,(h\theta/b)\cos[\omega t - (z/b)(2\pi-\theta)]}{\mathrm{ch}\,(2\pi h/b) - \cos(2\pi z/b)}. \qquad (2.2)$$

If $2\pi h/b \gg 1$, which is usually the case, this expression can be replaced with adequate accuracy by the formula

$$H(z, t) = \frac{4\pi J_0}{bc}\cos\left(\omega t + \frac{\theta z}{b}\right) e^{-\theta h/b}. \qquad (2.2a)$$

Hence it is clear that when θ is neither zero nor π, we have a traveling wave and the field strength increases from the surface of the plasma to the periphery. The phase difference is $\theta = 2\pi/m$ where m is the number of phases and $\theta/b = \kappa = 2\pi/\lambda$, where λ is the space wavelength. Hence it follows that the wave propagates in the z direction with the phase velocity $v_{ph} = \omega/\kappa$ and the depth of the magnetic well increases as the ratio of the wavelength λ to the distance h decreases.

The magnetic pressure on the plasma, averaged over the coordinates and time, is

$$\overline{P}_M = \frac{1}{2\pi b}\int_{-\pi}^{\pi}\left(\int_{-\pi}^{\pi} H^2(z, t)\, d\omega t\right) dz = \pi\left(\frac{J_0}{bc}\right)^2 \frac{\mathrm{ch}\,[(2\pi h/b) - 2\kappa h]}{\mathrm{sh}\,(2\pi h/b)}; \qquad (2.3)$$

where we have used (2.2). We now substitute the approximate formula given by (2.2a) and obtain

$$\overline{P}_M = \pi\left(\frac{J_0}{bc}\right)^2 e^{-2\varkappa h}. \qquad (2.3a)$$

This average pressure should balance the gas dynamic pressure in the plasma and, therefore,

$$\pi\left(\frac{J_0}{bc}\right)^2 \frac{\operatorname{ch}[(2\pi/bc)-2\varkappa h]}{\operatorname{sh}(2\pi h/b)} \approx \pi\left(\frac{J_0}{bc}\right)^2 e^{-2\varkappa h} = nkT, \qquad (2.4)$$

where n is the plasma density, T is the temperature, and k is Boltzmann's constant.

Since the true pressure on the plasma at any given time is different from the average at each point, the plasma surface exhibits perturbations. The size of the perturbation, $\Delta(z, t)$ (see Fig. 1), must be small in comparison with the distance h between the plasma surface and the conductors. The maximum distortion of the plasma surface was calculated in [2] for two cases: the three-phase (m = 3) and six-phase (m = 6) systems and the results are, respectively,

$$\left(\frac{\Delta}{h}\right)_{\rm extr} = 0.091 \text{ and } \left(\frac{\Delta}{h}\right)_{\rm extr} = 0.15.$$

These two values may be regarded as acceptable. The curvature of the "magnetic well" for m = 3 is greater than for m = 6, and this provides more favorable conditions for equilibrium and thermal insulation of the plasma. The three-phase system was therefore chosen in the experiment although the magnetic-field energy in the region between the plasma and the conductors for given magnetic pressure on the plasma surface should be somewhat higher than for m = 6.

The above results enable us to choose the currents in the coils, the number of phases, and the ratios of the basic parameters including, in particular, h/b.

All that remains is to find the criterion for the frequency, which is also determined from the condition that the surface distortion must be small at different instants of time, and not at different points for a fixed instant of time, as was done above. The corresponding calculations are also given in [2] and lead to the condition

$$\omega > \frac{v_{Ti}}{4h}, \qquad (2.5)$$

which determines the minimum frequency, i.e., the required frequency band lies in the region of a few megahertz.

The corresponding wavelength of the electromagnetic wave amounts to a few hundred meters, which is much greater than all the geometric dimensions of the system.

We have thus determined all the basic parameters of devices of this kind, subject to the simplifying assumptions adopted in the calculations.

More detailed analysis shows that there are several important effects which have not been taken into account above. These include, above all, nonlinear effects which appear as a result of the drift of electrons dragged into motion by the traveling wave. Electron drift leads to the appearance of a current in the z direction, and this current has a constant component.

In order to investigate this phenomenon, we shall write out the complete set of equations describing electromagnetic processes in plasma exposed to a traveling magnetic field. We must introduce the finite conductivity σ of plasma, and investigate the field penetrating the plasma. The characteristic length representing the depth of penetration of the field into the plasma is the skin depth δ. The significance of the skin layer in the present problem will be evaluated below. Since, usually, the depth of the skin layer is also much less than all the other characteristic dimensions, we can confine our attention to the plane problem.

The basic equations written in the usual notation are the generalized Ohm's law

$$\mathbf{j} = \sigma\mathbf{E} + \frac{e\tau}{mc}[\mathbf{j}, \times \mathbf{H}] \qquad (2.6)$$

and the Maxwell equations (without displacement currents)

$$\operatorname{rot}\mathbf{H} = \frac{4\pi\mathbf{j}}{c}, \quad \operatorname{rot}\mathbf{E} = -\frac{1}{c}\frac{\partial\mathbf{H}}{\partial t}, \quad \operatorname{div}\mathbf{H} = 0. \qquad (2.7)$$

We shall suppose that the plasma temperature and density are everywhere constant, so that $\sigma = e^2 n\tau/m$ does not vary throughout the plasma. This may not be a fully justified assumption but must be introduced since it provides us with the only possibility of solving the problem.

The solution obtained for this idealized problem leads to a qualitative picture of the phenomenon.

The situation is, however, complicated still further by the following fact. We are using the usual, i.e., collisional, conductivity. In this connection, it is usually assumed that the skin layer depth is much greater than the mean free path for electrons undergoing Coulomb collisions. In reality, this condition is not satisfied in the most interesting cases.

The Coulomb mean free path is given by (in cm)

$$l = \frac{3\cdot 10^{13}}{n_e L_C} T^2 \text{ cm};$$

where T is in electron-volts and L_C is the Coulomb logarithm which we shall assume to be approximately 10. If the temperature is 100 eV and $n_e = 10^{15}$ cm^{-3}, we have $l \approx 30$ cm, whereas the skin layer depth is less than 1 cm.

The existing theory of the anomalous skin effect is not valid in this case because it is based on the assumption that the pressure due to the high-frequency field produces small density perturbations. In our case, the field pressure completely balances the plasma pressure, so that the distance over which the plasma density changes from zero to its maximum value is comparable with the skin-layer depth.

In a recent paper, Halland and Gerwin [5] considered the penetration of the high-frequency field into plasma, when the pressure of this field was sufficient for the confinement of dense hot plasma, on the assumption that collisions could be neglected.

Their results (obtained with the aid of a computer) can be briefly summarized as follows. The skin layer depth for a relatively broad range of the corresponding parameters is always equal to a few units of c/ω_0 [the plasma frequency is $\omega_0 = (4\pi e^2 n_e/m)^{1/2}$], or

$$\delta = k\frac{5\cdot 10^5}{\sqrt{n_e}}, \qquad (2.8)$$

where the numerical coefficient is $k \approx 2-3$.

The second result is that the phase of the wave does not change in the course of penetration into the plasma and remains constant at the value corresponding to its value in vacuum.

It is clear that the quantity δ given by (2.8), subject to the condition that $H^2/8\pi = n_e kT$, reduces to

$$\delta = k\frac{v_{Te}}{\omega_H} = kr_L, \qquad (2.9)$$

where v_{Te} is the electron thermal velocity and r_L is the electron Larmor radius.

Consequently, the results reported in [5] mean that the skin layer depth amounts to a few electron Larmor radii, which correspond to the effective value of the high-frequency field on the plasma-vacuum boundary.

We shall see below that these conclusions reveal a certain conflict with the experimental results. The skin layer depth is found to exceed the prediction based on (2.8) and, at low densities the skin layer, tends to fall rather than increase. All this refers to experiments with the best parameters (high fields and low densities) which are of particular interest because the usual ideas are then known not to apply. Moreover, a strong change in the phase is observed as the field penetrates the skin layer. The phase shift reaches up to 100–140°. It is important to remember, however, that under the conditions of these experiments there was a considerable drift current in the plasma, and the field due to this current was comparable with the applied field. This may have had a substantial effect on the penetration of the high-frequency field into the plasma.

Let us now return to our basic problem. This is expressed by (2.6) and (2.7) which, under our adopted assumptions, form a closed system whose solution will yield all the quantities in which we are interested.

Since (2.6) is nonlinear and quite complicated, quantitative results can only be obtained by numerical integration of (2.6) and (2.7), subject to the corresponding boundary conditions. This problem has been solved on a computer and the results are published in [6-9]. The difference between [6,7] and [8,9] is in the boundary conditions and characteristic parameters. Thus, in [6,7] the field is specified on the plasma surface whereas in [8,9] the current in the coil is specified.

We shall not consider a largely qualitative analysis (see [10]) which will help us to elucidate the basic physical processes which take place in plasma under the above conditions.

Since the processes in which we are interested occur in a thin skin layer, we shall again confine our attention to the plane problem. All the y derivatives will therefore be assumed to be equal to zero. Moreover, we shall assume that the current component $\partial H_y/\partial z$, which is perpendicular to the plasma surface, is negligible and that $\partial H_x/\partial z$ is small in comparison with $\partial H_x/\partial x$. The ratio of these quantities is of the order of $(\kappa\delta)^2 \ll 1$ where δ is the skin layer depth which is much smaller than the space wavelength λ of the electromagnetic wave.

We shall next assume that as the wave penetrates the plasma the magnetic field of the traveling wave changes in the same way as for a linear system, i.e.,

$$H_z = H_{z0}e^{-x/\delta}\cos\left(\omega t + \varkappa z - \frac{sx}{\delta}\right), \qquad (2.10)$$

where H_{z0} is the amplitude of the high-frequency field on the surface of the plasma and s is a coefficient which governs the phase shift of the wave and must be found experimentally.

Under the above assumptions, the equations given by (2.6) and (2.7) can readily be solved and the result is

$$j_y = \frac{\sigma E_y}{1 + (e\tau H_x/mc)^2}. \qquad (2.11)$$

This is the generalized Ohm's law which can be written in the more usual form by introducing the effective conductivity

$$\sigma_{\text{eff}} = \frac{\sigma}{1 + (e\tau H_x/mc)^2}. \qquad (2.12)$$

The value of H_x can readily be found from the condition that $\operatorname{div} H = 0$, and for H_z we can use (2.10).

In addition to the high-frequency current j_y, the solution of (2.6) and (2.7) enables us to find the drift current in the z direction, whose density is given by

$$j_z = -\frac{en_e v_{\text{ph}}(e\tau H_x/mc)^2}{1 + (e\tau H_x/mc)^2}. \qquad (2.13)$$

It is clear that this current is determined by the drift of electrons in the traveling wave field, and that the electron velocity is determined by the phase velocity of the wave and cannot exceed it. Moreover, it follows from (2.13) that the drift current contains a time-independent component. The density of this constant component is

$$\bar{j}_z = \frac{1}{2\pi}\int_{-\pi}^{\pi} j_z\, d\omega t = -env_{\text{ph}}\left(1 - \frac{1}{\sqrt{1 + q^2 e^{-\frac{2x}{\delta}}}}\right). \qquad (2.13a)$$

In this expression, we have introduced the dimensionless parameter

$$q = \frac{e\tau}{mc}\varkappa\delta H_{x0}.$$

The solution obtained for the plane wave can be generalized to more realistic conditions, namely, a cylinder of radius a whose thin walls carry currents. Integrating (2.13a) with respect to x, and assuming that the current flows over the perimeter of the cylinder $2\pi a$, we obtain the constant component of the drift current (which can be measured) in the form

$$J_z = 2\pi a\int_0^\infty \bar{j}_z\, dx = -\pi a e n_e v_{\text{ph}}\delta \ln\frac{q^2(1+\sqrt{1+q^2})}{4(\sqrt{1+q^2}-1)}. \qquad (2.14)$$

Assuming that this current produces a field H_φ around the cylinder, we can find the ratio of the field due to the constant component of the drift current on the plasma vacuum boundary to the amplitude H_{z0} of the high-frequency field:

$$\frac{H_\varphi}{H_{z0}} = \frac{\delta_\text{K}^2}{\delta^2}\frac{1}{\sqrt{1+s^2}\,q}\ln\frac{q^2(1+\sqrt{1+q^2})}{4(\sqrt{1+q^2}-1)}, \qquad (2.15)$$

$$\delta_\text{C}^2 = \frac{c^2}{4\pi\sigma\omega}.$$

When looked upon as a function of q, this ratio has a well-defined maximum at $q \approx 2.4$, and the function of q on the right-hand side of (2.15) then has the value of ≈ 0.49.

In the above discussion, the skin-layer depth δ is assumed given. The value of δ is found from the self-consistent problem described by (2.6)–(2.7), the numerical solution of which [6-9] has led to the following results. The problem is solved on the assumption that the conductivity is governed by the usual Coulomb collision mechanism. It then turns out that the skin-layer depth δ under the above conditions is increased by a factor of about 2, as compared with the usual depth, and further increase in the corresponding parameters leads to saturation in the value of the skin-layer depth. Numerical calculations show also that the ratio H_φ/H_{z0} is bounded, and that its maximum value is approximately 0.8, which corresponds to $q \approx 3.5$.

It thus turns out that even in this system which, at first sight, looks like an ideal trap, there is a "special" direction which is connected with the appearance of the drift current whose field is comparable with the applied field for the range of parameters in which we are interested. This leads to at least three important consequences. It is well known that the plasma ring cannot be in equilibrium on its own. It is subject to the electrodynamic repulsive force

$$F_1 = -\frac{1}{2c^2}\frac{\partial}{\partial R} LJ_z^2 = -\frac{2\pi}{c^2} J_z^2 \left(\ln\frac{8R}{a} - 1\right), \qquad (2.15')$$

where J_z is the drift current, and R and a are, respectively, the major and minor radius of the plasma ring.

It is readily shown that this force cannot, in practice, be compensated by the pressure associated with the external high-frequency field. It follows that this system, which might be imagined to be the ideal trap, is, in reality, an unstable configuration. However, this difficulty can be overcome in two ways. First, the current ring can be "suspended" in a transverse field.[11] If the transverse field is denoted by $H_\perp$, then the force associated with this field, which is directed toward the center of the torus, is $F_2 = 2\pi R J_z H_\perp/c$. The equilibrium condition then reduces to

$$H_\perp = \frac{J_z}{Rc}\left(\ln\frac{8R}{a} - 1\right). \qquad (2.16)$$

It is clear that this field is small and should not lead to additional difficulties. Moreover, since the process is stationary (the current is established in a short time), this type of equilibrium or, more precisely, the control of the position of the plasma ring relative to the axis of the discharge chamber, can be maintained without much difficulty throughout the course of the process. This has, in fact, been achieved in practice. The plasma ring suspended in this way has definite advantages, since the field due to the drift current produces additional pressure on the surface of the plasma and this adds to the pressure associated with the high frequency field and smooths out effects connected with the presence of rapidly varying fields.

The drift current, and the complications associated with it, can be almost totally eliminated in practice. This is achieved with the aid of a system of two traveling waves moving in opposite directions with roughly equal amplitudes and phase velocities. It is clear that these waves must have different frequencies and different space wavelengths because the condition $v_{ph1} = v_{ph2}$ leads to $\omega_1/\omega_2 = \lambda_2/\lambda_1$. In principle, this system exhibits all the properties of the ideal trap, since it does not contain special directions. A system with waves traveling in opposite directions has been constructed and investigated. The results will be reviewed below.

However, the development of this system leads to additional technological difficulties. Therefore, because experiments on the suspended-ring system lead to sufficiently satisfactory results, this system must be regarded as offering the best solution for the moment. It is also simpler to realize in practice than the system with opposing waves. There is a further fact associated with the presence of the drift current, namely, the possible appearance of MHD instabilities. It is well known that the plasma column containing a longitudinal current is unstable against perturbations with $m = 0$ (sausage type) and $m = 1$ (kink and screw type instabilities). In our case, the plasma column is under the action of a longitudinal high-frequency field for which the stability conditions are much more readily satisfied than for a column carrying a current in a stationary longitudinal field. This question will be discussed in Sec. 4 under the heading of dynamic stabilization. Here, we merely note that in the experiments performed so far, in which the equilibrium conditions were satisfied, no evidence has been found for the presence of instabilities, at least at a macroscopic scale.

The presence of the drift current in the plasma leads to additional heating which is connected with the large consumption of active power.

In fact, the power dissipated per unit volume of plasma is

$$P = \frac{j^2}{\sigma_{eff}}, \quad \sigma_{eff} = \frac{\sigma}{1 + (e\tau H_x/mc)},$$

and since $j_z = -(e\tau/mc) j_y H_x$, we have

$$P = \frac{j_y^2}{\sigma} + \frac{j_z^2}{\sigma}. \qquad (2.17)$$

Hence, it follows that σ_{eff} characterizes the high-frequency energy dissipation per unit volume of the plasma, plus the additional energy which must be expended to maintain the drift current and which is dissipated when this current passes through the plasma.

Whilst in systems with constant and slowly varying fields the production of hot plasma at the required temperature is a major problem for which an adequate solution has not yet been found, in the case of plasma confinement and stabilization by high-frequency fields, we have the opposite problem, i.e., "plasma overheating."

When (2.20) was derived, it was assumed that the energy was released in the region through which the current flowed, i.e., the skin layer. This is valid in the usual skin layer for which the skin layer depth is greater than the electron mean free path. In the case of the anomalous skin effect, and small magnetic fields, the skin layer depth is less than the mean free path but greater than the electron Larmor radius. Under these conditions, the electron acquires energy from the electric field in the skin layer, but loses it by collisions with other particles, mainly in the part of the plasma where there are practically no fields. In this case, (2.20) will not, of course, be valid.

In the case of interest to us, the electron Larmor radius is less than the skin layer depth. Therefore, an electron drawing energy from the electric field in the skin layer will drift in the magnetic field and remain in the region of sufficiently strong fields until, as a result of collisions with other particles, it escapes into the low-field region. Therefore, provided we exercise care, we can use (2.20) in this case. For the first term in this expression, we find, using cylindrical coordinates, that the total power dissipated per unit length in the plasma is

$$P = \frac{2\pi}{\sigma}\int_0^a j_\varphi r\, dr = \frac{H_{z0}^2}{8\delta} a\delta_C^2 \omega;$$

where we are assuming that $H_z = H_{z0}\exp[-(a-r)/\delta]$ and δ is the anomalous skin layer depth ($\delta \ll a$).

The energy liberated in the plasma per unit volume (per unit length) in a time τ is

$$W = \frac{H_{z0}^2 a \delta_C^2 \omega\tau}{8\delta}.$$

The thermal energy per unit volume is $nkt = H_{Z0}^2/8\pi$ (where H_{Z0} is always the root mean square effective field). On the other hand, by comparing these quantities, we can obtain a formula which will enable us to find the time for overheating. During this time, the plasma heats up (as a result of the application of the high-frequency field) to a temperature at which the gas pressure is balanced by the magnetic pressure. The corresponding condition is

$$\frac{H_{z0}^2}{8\pi} = \frac{H_{z0}^2}{8\pi}\frac{\omega\delta_C^2\tau_{ov}}{a\delta};$$

and hence

$$\tau_{ov} = \frac{a\delta}{\omega\delta_C^2} = \frac{2\pi a\sigma\delta}{c^2}. \qquad (2.18)$$

If we assume, following [5], that $\delta = kc/\omega_0$, where k is of the order of a few units, then

$$\tau_{ov} = \frac{2\pi ak\sigma}{\omega_0 c} = \frac{5\cdot 10^{-2} T_{(eV)}^{3/2} ak}{\sqrt{n}}. \qquad (2.18a)$$

As an example, consider the conditions which must be satisfied by a system of this kind, working as a thermonuclear reactor with a positive energy yield.

To ensure that the energy released during thermonuclear fusion in the DT reaction is greater than the energy input into the plasma, we must have

$$n\tau > 10^{14}, \qquad (2.19)$$

where τ is the lifetime of the particles in plasma (seconds) and n is the plasma density (cm^{-3}). The plasma ion temperature is then T = 10 keV.

We shall suppose that the lifetime is equal to the time for overheating, in which case (2.18a) yields (taking k = 3)

$\pi a^2 n = N = 3\cdot 10^{18}$ particles/cm so that $\tau_{ov} \approx 1$ sec

If we take a = 100 cm, then $H_{Z0} = 10^4$ Oe. These figures may appear difficult to achieve at present, but will probably become quite realistic in the near future.

However, we now encounter the basic difficulties mentioned above.

Equation (2.19) was derived on the assumption that the entire energy was expended in heating the plasma and that the energy losses in all the electrical supply systems were negligible. This requirement can be satisfied in practice when the current in the conductors supplying the system varies sufficiently slowly. For high-frequency devices, however, the energy losses in the supply system cannot be reduced to a sufficiently low level because of the well-defined skin effect in the conductors.

The energy losses in the circuit are given by

$$P_f = \frac{\omega\tau}{Q}\int_{V_f}\frac{H^2}{16\pi}dV,$$

where the integral is evaluated over the entire volume V_f occupied by the high-frequency field, Q is the Q-factor of the oscillatory system, and H is the maximum magnetic field. The thermal energy stored in the plasma is

$$P_p = \int_{V_p} nkT\, dV,$$

where the integral is evaluated over the volume V_p occupied by the plasma.

The pressure balance condition is $H_0^2/16\pi = nkT$ (H_0 is the field on the surface of the plasma, whilst nkT is taken on the axis where the pressure is a maximum) and, therefore, these two expressions give us a condition which ensures that the losses in the supply circuit are less than the thermal energy in the plasma. This condition can be reduced to

$$Q > \omega\tau K_f\frac{V_f}{V_p}; \qquad (2.20)$$

where K_f is the "form factor" which characterizes the variation of the magnetic pressure with distance from the surface of the plasma in the peripheral direction. If this field increases toward the periphery, then K_f is greater than unity and increases with increasing field curvature.

When K_f is not too large, and the field and plasma volumes are usually of the same order, (2.20) can be written, for our purposes, in the form

$$Q > \omega\tau.$$

From (1.1) we have $\omega \approx 10^8$ and, if we suppose that $\tau \approx 1$ sec, we find that $Q > 10^8$, whereas the usual Q factor of high-frequency systems in this frequency region is ~150–200. When the frequency is increased, the situation becomes worse still because Q increases with frequency only as $\omega^{1/2}$.

Superconducting coils are being increasingly used in modern systems with constant and quasiconstant fields, designed for investigations into the question of controlled thermonuclear fusion. It is now possible to construct superconductors whose properties are unaffected by sufficiently strong fields (up to some tens of kOe). However, these are type-II superconductors which retain their properties only for stationary currents. The characteristic time for the change in the currents cannot then be less than of the order of the second. It is clear that such superconductors cannot be used in high-frequency devices. It follows that high-frequency circuits with the necessary value of Q cannot as yet be constructed. There are, however, some grounds for moderate optimism. It is well known that, some years ago, Little put forward the idea that it may be possible to produce superconductors with very high critical temperature. One would hope that a solution free from the disadvantages exhibited by type II superconductors will be found in the not too distant future. The confinement and stabilization of plasma by high-frequency fields for the purposes of thermonuclear fusion will become technologically possible when high-Q circuits for fields of some tens of kOe are developed. Until this happens, this approach to thermonuclear fusion can only be regarded as a problem in physics and not a technological possibility. Nevertheless, the interaction between strong high-frequency fields and plasmas is in itself exceedingly interesting and may turn out to be very useful in other branches of physics.

We believe that such studies should be continued and developed in very different directions in order to prepare the ground for future technological solutions.

After this brief excursion into futurology, we must now return to reality and try to compare existing results with the various relationships given above. In studies performed on the "Volna" installation, the typical values were $\delta = 0.5$ cm, $T_e = 25$ eV, and $a = 4$ cm. From (2.18) we then have

$$\tau_{ov} = \frac{2\pi a\delta\sigma}{c^2} \approx 25 \cdot 10^{-6} \text{ sec}$$

and since, according to (2.17), losses due to the drift current are approximately of the same order, the true time for overheating must be reduced by a factor of two, in which case $\tau \approx 12$ μsec, which is quite close to the observed value.

Therefore, to explain the phenomena observed in experiments performed so far, there is no need to introduce any additional energy dissipation mechanisms into the plasma. So far, the observed picture lies wholly within the framework of classical ideas on the collisional mechanism of heating. It must not, of course, be forgotten that the skin layer depth used in the above calculations is taken from experiment and is not the usual skin layer depth.

3. EXPERIMENTAL STUDIES OF THE "TRAVELING WAVE" SYSTEM

These experiments were performed on several installations. The most complete and reliable results were obtained for the Delta-2 and Volna installations. These results were published in[10,12-21].

We must first note that the development of the high-frequency system, namely, the supply generator and the matching of it to the very unusual load, lead to very considerable difficulties. A number of specific problems has to be solved in the development of such systems: the currents producing the high-frequency magnetic fields must be kept relatively constant while the plasma conductivity undergoes a very considerable change, the high-frequency oscillations must be excited at a frequency which follows sufficiently closely the natural frequency of the circuit loaded with the plasma, the traveling magnetic field must be excited in the coil and the direction of propagation of the traveling wave must be stabilized, a high-frequency system must be developed with power output of several tens of megawatts for pulse lengths of the order of milliseconds, and the output power of many simultaneously operating vacuum-tube generators must be added. We shall not consider in detail the various engineering problems and their solutions, many of which are highly original. This has been a challenge for the skill and originality of the engineers, led by I. Kh. Nevyazhskiĭ, who developed the system. A few words must be introduced in this connection.

The maximum power per pulse which can be derived from a single tube is approximately 5 MW. The required active power is

$$P = \frac{\hat{H}^2}{16\pi}\frac{\omega V}{Q}$$

where V is the working volume. The experiments were carried out on a toroidal chamber with $R = 40$ cm and $r = 5$ cm, where r is the radius of the current-carrying coil which defines the field volume. At first, we used a quartz chamber and, subsequently, a ceramic one. Under these conditions, $V = 2\pi^2 Rr^2 = 2 \times 10^4$ cm^3 and $\omega \approx 10^7$. The field amplitude, for which experiments of basic physical interest can be performed, is $H \approx 10^3$ Oe, in which case

$$P = \frac{10^6 \cdot 10^7 \cdot 2 \cdot 10^4}{50Q} \cdot 10^{-7} = \frac{400}{Q} \text{ MW}.$$

Since the efficiency of the system is ≈ 0.3, and if we assume that $Q = 20$, which is close to the observed value, we find that the required power output of the vacuum-tube generator should be about 60 MW. The vacuum-tube generator therefore consisted of twelve 5-MW tubes. The power outputs of these tubes had to be combined in order to produce the traveling wave. This was done with the aid of a three-phase system with a phase shift of 120°. The system is based on coupled circuits with the secondary detuned on open circuit, so that near-resonant conditions are produced when the circuit is loaded in the presence of the plasma. The vacuum-tube generators are self-excited but this substantially reduces the efficiency of the system due to high grid currents. However, this makes the system flexible and insensitive to changes in circuit parameters. The traveling field is thus produced in an artificial LC line. The helix which forms the inductance of the line is closed into a ring, forming a torus containing the discharge chamber. A capacitance is connected between each turn of the helix and the common grounding rail. The size of the chamber was indicated above. The number of turns is 180 and the pitch of the helix is about 12 mm. Ceramic capacitors of 470 pF each are used in the artificial line. The total capacitance of the system is 0.367 μF. From the spectrum of resonance frequencies of the system, the working frequency was chosen to be $f = 2.5$ MHz, for which 12 waves fit into the perimeter of the torus. The space wavelength is therefore $\lambda = 18$ cm. This wavelength and, of course, the working frequency can be varied within very broad limits. In the latest measurements, the number of waves was varied from 4 to 16. The experiments described here refer to a system of 12 waves. The field near the plasma surface then rises sufficiently rapidly in the peripheral direction and this improves the plasma stability conditions. Field distortions, due to the fact that the system is toroidal and the pitch of the helix finite, are quite small. The length of the high-frequency pulse can be raised up to 500 μsec.

To be able to carry out the experimental work at sufficiently low initial gas pressure in the chamber (less than 10^{-3} Torr), we used an additional generator for preliminary ionization with an output power of about 100 kW at 13 MHz. The pulse from this generator preceded the main working pulse.

The first cycle of measurements was concerned with the distribution of the magnetic fields, namely, the applied high-frequency frequency field and the constant component of the field due to the drift current. These measurements were carried out with miniaturized magnetic probes. The diameter of these probes did not exceed 1–2 mm, i.e., it was much smaller than the characteristic dimensions of the system. The magnetic probes were placed in quartz tubes (outer diameter of 3–4 mm) placed across the discharge chamber in horizontal and vertical directions. The probes could be moved along these tubes so that the magnetic fields could be found at at all the required points.

The distribution of the high frequency and quasi-stationary fields in the plasma enabled us to elucidate the interaction between the field and the plasma, depending on the initial gas pressure and the amplitude of the applied magnetic field. We reproduce the characteristic distribution curves for the high-frequency field in the plasma over the cross section of the discharge chamber (z component) under different conditions, together with the corresponding oscillograms of the high-frequency field amplitude and the drift current produced by the traveling wave.

The form of the high-frequency field distribution suggests the existence of two characteristic regimes, depending on the values of the above parameters. The first (Fig. 2a) corresponds to initial gas pressures in the range $0.01 < p_0 < 0.1$ Torr. When the emplitude of the applied field is up to 500 Oe, there is a monotonic fall in the field between the periphery and the center of the discharge chamber. Here, we can see a well-defined skin layer. The skin layer depth varies somewhat, depending on the field: it decreases with increasing field. No explicit dependence on time beyond a certain moment is observed. Even under these relatively weak conditions, the skin layer depth turns out to be several times greater than the ordinary skin depth. At the walls of the chamber, where the field is a maximum, the effective conductivity is low but inside the plasma, where there is practically no field, the electron temperature measured by the double probe amounts to 10–12 eV. Figure 2b shows the drift-current oscillograms under these conditions, and Fig. 2c shows a typical oscillogram of the hf field envelope.

When the initial pressure is reduced to about 2×10^{-3} Torr, the hf field distribution undergoes a substantial change and one observes a relatively well defined time dependence. Figure 3a shows the distribution function at different instants of time for a maximum field of 600 Oe. It is clear that, eventually, the plasma partially departs from the inner wall of the chamber. This is accompanied by a reduction in the skin-layer depth, and the drift current increases rapidly (Fig. 3b).

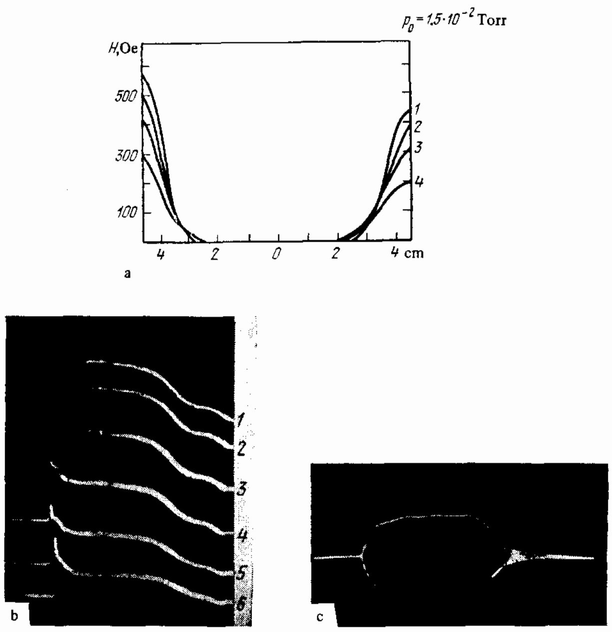


FIG. 2. (a) Distribution of the hf field amplitude $\tilde{H}_z$ over the cross section of the discharge chamber at a fixed instant of time for different anode voltages U_a of the generator tubes (all the measurements are referred to the axis of the chamber cross section; the initial pressure is $p_0 = 1.5 \times 10^{-2}$ Torr, U_a (kV) = 26 (1), 22 (2), 18 (3), and 14 (4); (b) oscillogram showing the drift current for different U_a with $p_0 = 1.5 \times 10^{-2}$ Torr [U_a (kV) = 26 (1), 22 (2), 20 (3), 18 (4), 16 (5), and 14 (6)]; (c) typical oscillogram of the envelope of the hf field at the wall of the discharge chamber for relatively high pressure $10^{-2} < p < 0.1$ Torr (pulse length 500 usec).

The $\tilde{H}_z$ oscillograms exhibit weak oscillations. Further reduction in the initial pressure leads to a clear separation of the plasma from the inner wall (Fig. 4a). The curves given in this figure correspond to a fixed instant of time somewhere near the center of the applied pulse. Oscillograms of the hf pulse envelope and the drift current (Fig. 4b) show rapid oscillations with large amplitudes which, at first sight, may appear as a sign of instability. More detailed analysis shows, however, that this is not an instability and that the oscillations are connected with loss of equilibrium due to the fact that the drift current rises sharply under these conditions. It was shown above that the electrodynamic repulsive forces acting on the current-carrying plasma ring cannot be compensated, when this current is large, by the pressure associated with the high-frequency field which decreases toward the center, and that this system cannot be in equilibrium.

Plasma which interacts strongly with the walls under-

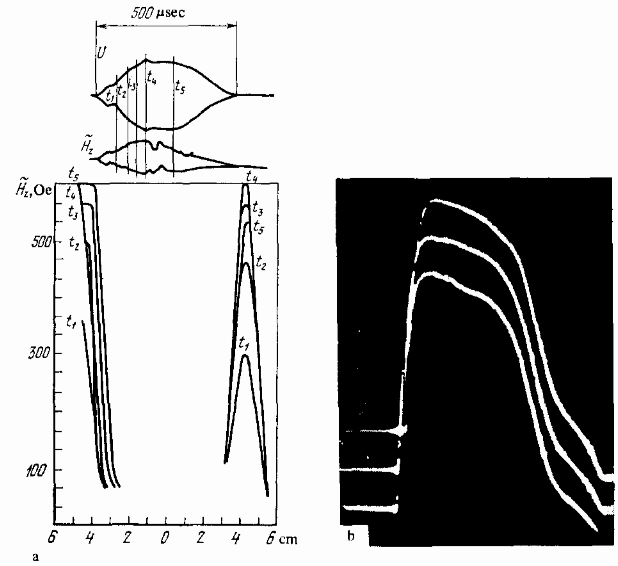


FIG. 3. (a) Distribution of the amplitude of the hf field $\tilde{H}_z$ over the cross section of the discharge chamber at different times for an initial pressure $p_0 = 2 \times 10^{-3}$ and anode voltage $U_a = 26$ kV (oscillograms of the envelope of the voltage across the circuit and of the hf field at the chamber wall are shown at the top of the figure; we also show the time scale and instants corresponding to points on the distribution curve); (b) oscillograms showing the drift current at anode voltage $U_a = 26$ kV for the upper oscillogram, 22 kV for the middle, and 20 kV for the bottom oscillogram (same scale as in Fig. 2b; maximum drift current J = 3 kA).

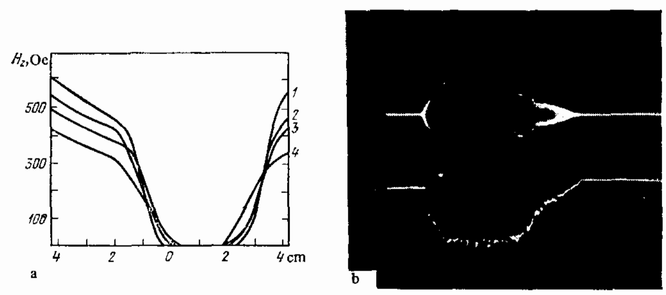


FIG. 4. (a) Distribution of the amplitude of the hf field $\tilde{H}_z$ over the cross section of the discharge chamber at a given time for U_a (kV) = 22 (1), 20 (2), 18 (3), and 16 (4) (initial pressure $p_0 = 8 \times 10^{-4}$ Torr); (b) oscillograms of the envelope of the hf field at the wall of the discharge chamber (upper trace) and the drift current (lower trace) for $p_0 = 8 \times 10^{14}$ Torr and $U_a = 22$ kV.

goes cooling, the drift current falls, the equilibrium is re-established, the plasma moves away from the walls and is heated, the drift current increases, the equilibrium is violated, and so on. It is therefore necessary either to compensate the electrodynamic repulsive forces acting on the plasma ring by the force of interaction between the drift current and the transverse magnetic field, or to eliminate the drift current by using traveling waves propagating in opposite directions.

These possibilities were investigated on the Volna installation in a natural continuation of the experiments performed previously on Delta-2.

Both in its design and in the high-frequency supplies, the Volna system is similar to Delta-2. The essential point is that the Volna generator was divided into two groups of three-phase generators, and the circuit was modified so that the following operating conditions could be achieved:

1) Separate operation of each of the generators with one group working at 1.9 MHz or 1.6 MHz, and the other group at 2.8 or 3.8 MHz.

These frequencies correspond to 4, 8, 12, and 16 waves fitting into the coil.

2) Simultaneous operation of both generators at the frequencies indicated above with the pulse from one generator shifted in time relative to the pulse from the other, or the simultaneous generation of the two pulses when the two opposite traveling waves are employed.

3) Combined operation of the two generators in the form of a single six-phase or three-phase system at one of the above frequencies in order to achieve maximum output power.

In the Volna installation, a transverse magnetic field was used to compensate the electrodynamic repulsive forces on the plasma ring due to the drift current when only one traveling wave was employed. This transverse field was produced by placing two or four circular turns on the outer perimeter of the toroidal discharge chamber, which were supplied by the modulator feeding the vacuum-tube generator.

Figure 5 shows the measured and calculated distributions of the amplitude of the longitudinal component of the hf field for the above numbers of waves.

It is clear that, as the number of waves increases, the field increases in the direction of the wall more rapidly and this, of course, improves the stability conditions. However, there is an attendant increase in the toroidal inhomogeneity, and the possible repulsion of the plasma in the direction of the outer wall becomes increasingly dangerous. This effect is compensated by the interaction between the drift current and the transverse field. It is known [11] that, under these conditions, the equilibrium is stable in the horizontal plane for the plasma ring as a whole. However, very stringent conditions must be satisfied by the transverse-field configuration if the ring is to retain its stability in the vertical plane as well. In our case, the presence of a high-frequency field which is symmetric in the vertical plane, and has a well-defined "well," ensures stability in this direction.

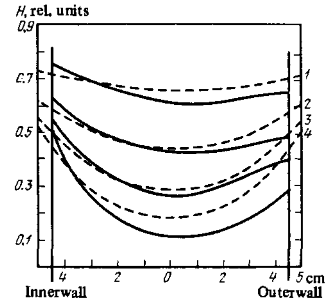
This method of controlling the equilibrium position of the plasma ring tends to conflict with the original idea because we are now producing a special direction, i.e., we no longer have an ideal trap in the sense indicated above.

Nevertheless, the results show that, in this case, one achieves not only equilibrium but also ring stability, at least against MHD-type instabilities.

Figure 6 shows oscillograms of the hf field recorded by a probe displaced in the horizontal direction over the cross section of the discharge chamber at a given time. It is clear that, depending on the ratio of the drift current to the transverse field, one can control the position of the plasma ring so that it remains separated from the walls at all points.

The top-left oscillogram shows the field distribution in the absence of compensation. The current in the coils producing the transverse field is $I = 0$. The middle oscillogram, which corresponds to $I = 500$ A, shows clear separation from both walls. This is, in fact, the optimum state. The bottom oscillogram shows overcompensation, i.e., the plasma presses against the inner wall of the chamber. Here, the current is $I = 1000$ A.

FIG. 5. Distribution of the amplitude of the hf field over the cross section of the working chamber for different frequencies (in the absence of a plasma): 1) wave No. 4, f = 0.9 MHz; 2) wave No. 8, f = 1.6 MHz; 3) wave No. 12, f = 2.6 MHz; 4) wave No. 16, f = 3.6 MHz. Dashed curves - theoretical; solid curves - measurements on a real system (in rel. units for equal values of the current in the exciting coil).



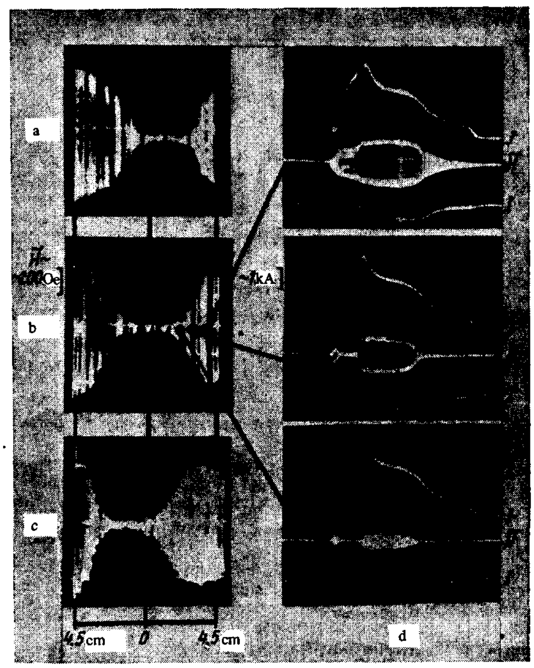


FIG. 6. Oscillograms of the hf field distribution: (a) without compensation, I = 0, plasma pressing against outer wall; (b) optimum compensation, I = 500 A; (c) overcompensation, I = 1000 A. plasma pressing against inner wall; (d) oscillograms of the drift current J, hf field $\tilde{H}$, and compensating current I corresponding to optimum compensation (oscillograms b). Oscillograms of the hf field $\tilde{H}$ were obtained for three positions of the probes, as shown in the figure. The scales of current, field, and time are indicated ($p_0 = 5 \times 10^{-3}$ Torr, 100 usec/cm).

On the right of the figure, we give oscillograms of the drift current J_2, the high-frequency field $\tilde{H}$, and the current producing the control field as functions of time.

The $\tilde{H}$ oscillograms on the right correspond to the three positions of the probe indicated in the figure. It is clear from these oscillograms that there is a time interval during which the probe records practically no field in a zone well away from the outer and, even more so, from the inner walls. This time we take to be the separation time.

This state of separation exists for 20 to 25 μsec, after which the plasma expands and reaches the chamber walls. When four space waves fit into the periphery of the torus, it is possible to increase the hf field amplitude up to about 1000 Oe. Under these conditions, the electron temperature measured by the double probe was found to be 60–70 eV.

The oscillograms in Fig. 6 refer to this situation. As noted above, the time of existence of the plasma ring is restricted by the "overheating time," whose magnitude is close to the observed value. Under these conditions, the Q factor of the circuit is found to be greater by a factor of two as compared with the values measured under different conditions. To increase the plasma temperature and its lifetime in the separated state still further, one must increase the output power of the generator.

The stability of this suspended plasma ring will be discussed below in Sec. 4 (dynamic stabilization). In this particular experiment, there were no signs of instability and no oscillations were detected, at least in the megahertz band and below.

We may conclude from this that the suspended ring in the traveling wave presents us with a promising system. It would be desirable to carry out further investigations in this area with improved supply-system parameters.

The variant in which two waves traveling in opposite directions are employed enables us to reduce the drift current, and the effects associated with the special direction can be reduced to a minimum. The real question is, of course, not which variant is closer to the ideal trap, but which method will enable us to produce plasma with the best parameters (temperature and density) for a given generator and with the equilibrium and stability conditions properly satisfied.

In this sense, the Volna system with the suspended ring offers the best solution. However, it is possible that, in the future, when the generator power output will not be subject to such serious restrictions, the system with the two waves traveling in opposite directions may become preferable.

Figure 7 shows oscillograms of the hf pulse envelope and the drift current which quite clearly show that oscillations connected with the loss of equilibrium due to the presence of the drift current when a single wave is employed (Figs. 7a and b) are suppressed when this current is compensated by the two opposing waves (Fig. 7c). It is clear from the oscillograms and the envelopes of the hf field that the system then enters a nonoscillatory state and is stable (Fig. 7d). It is also clear that there is an attendant substantial increase (by a factor of nearly 2) in the hf field, which suggests an increase in the Q of the plasma, i.e., a reduction in the resistive losses. This situation can be achieved and used in practice.

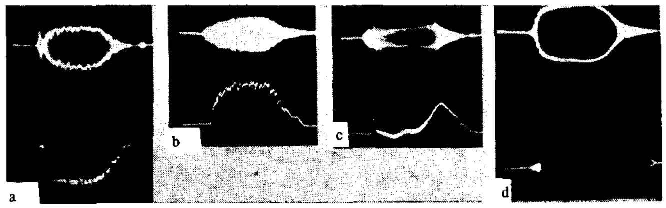


FIG. 7. Oscillograms of the envelope of the hf field (top) and drift current (bottom) during the operation of the Volna system with $p_0 = 3 \times 10^{-4}$ Torr. (a) Operation with one wave (8 waves fitting into the perimeter of the torus); $\tilde{H}$ = 400 Oe, maximum value of drift current J = 4 kA, frequency 1.8 MHz; (b) ditto for twelve waves directed in opposite direction; the fields and currents are roughly the same and the frequency is 2.8 MHz; (c) envelope of the hf field for two waves traveling in opposite directions (eight and twelve) (upper trace) and oscillogram of the drift current for two opposing waves (lower trace); (d) oscillograms of the hf field envelopes recorded with probes in the skin layer for 2.8 MHz (upper trace) and 1.8 MHz (lower trace); the oscillograms were recorded using filters tuned to these frequencies.

We have thus achieved our basic aim, i.e., to produce hot plasma, fully separated from the walls, with the thermal insulation conditions satisfied.

The lifetime of this separated state is unambiguously determined by the overheating time. There is some evidence that, as the generator output power increases, this time will also increase.

It was indicated above that there is at present no satisfactory theory of the penetration of the hf field into plasma under the conditions approaching the situation in traveling-wave systems.

To develop this theory, we must know the phase of the hf field as a function of penetration into the plasma. To obtain this information, we have carried out phase measurements as follows. The signal from the magnetic probe coil is fed along one channel through a bandpass filter to the plates of an oscillograph and through another channel to a phase-shift meter. The latter generates a pulse whose amplitude is proportional to the phase difference between the probe signal and a reference signal from a fixed loop located near the circuit. Figure 8 shows the phase shift of the hf field at approximately the center of the pulse.

These measurements were performed for the entire range of the parameters. It was found that the phase shift across the skin-layer depth was between $\pi/2$ and π, which is substantially greater than the ordinary shift and is certainly not consistent with the corresponding theories of the anomalous shift.

This concludes our review of experimental studies of systems of this type. We hope that the above data provide a clear enough idea of the difficulties and possibilities of these systems, and that they may stimulate further searches for improved variants of systems in which the plasma is confined by high-frequency fields.

4. DYNAMIC STABILIZATION: FUNDAMENTALS OF THE THEORY

It was noted above that the phrase "dynamic stabilization" was introduced when the stabilization of plasma by high-frequency fields was first suggested. The confinement of the plasma, i.e., the introduction of conditions ensuring the equilibrium of gas-kinetic and magnetic pressures, is then produced by constant or relatively slowly-varying magnetic fields. The high-frequency fields are used only for suppressing MHD-type instabilities. Here, too, we restrict our review to systems with waves in the megahertz range, in which the electromagnetic fields may be looked upon as quasi-stationary.

The first proposal for dynamic stabilization was concerned with the stabilization of the $m = 1$ mode (by high-frequency multipole fields [22] in the case of kink and screw instabilities, and by a longitudinal *hf* field [23] in the case of flute instabilities).

The basic idea of the first method of stabilization can be explained by considering a Z-pinch and a multipole with current directions as shown in Fig. 9. In this case, the pinch is stable as a whole during its motion along the horizontal axis and unstable along the vertical axis. If after a certain sufficiently short interval of time the currents in the multipoles and in the pinch are reversed, the pinch will be stable along the vertical axis and unstable along the horizontal axis.

Therefore, when alternating currents are present, there is a continual replacement of conditions favorable for instability by those unfavorable for instability, and vice versa. If the frequency at which these changes take place is sufficiently high, the pinch will be stable, on the average, against kink deformations, beginning with a certain value of the wavelength. The greater the kink wavelength, the easier is it to satisfy the stability conditions for this method. It is well known that when the pinch is stabilized by a longitudinal magnetic field the situation is reversed, i.e., as the kind wavelength increases, it becomes more difficult to satisfy the stability conditions.

The central idea of all the suggestions put forward so far in connection with the dynamic stabilization of plasma is that the spatial configuration of the magnetic field must change in time sufficiently rapidly in order to prevent the development of instability in a particular direction. At each given time, there are directions in which the system is unstable. However, it is possible to find certain definite conditions under which the system will retain stability in the presence of dynamic stabilizing

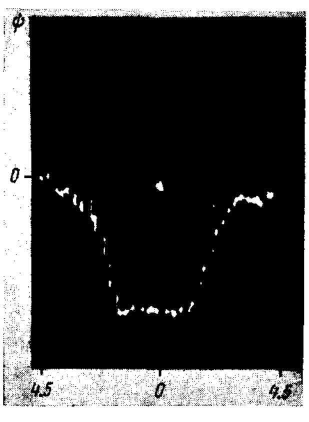


FIG. 8. Phase shift of the high-frequency field as a function of penetration into the plasma. The shift is measured relative to the phase at the wall of the discharge chamber; 1 cm corresponds to 60°. The maximum phase shift for this state is about 140°.

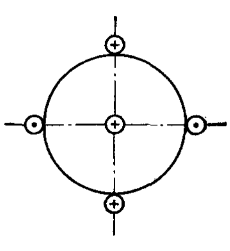


FIG. 9. Basic arrangement for dynamic stabilization.

fields which are much smaller than those necessary to maintain stability under stationary conditions.

Before we give the theoretical justification of the dynamic stabilization method, it will be useful to outline the elements of the theory of stabilization of a stationary Z-pinch by a longitudinal stationary magnetic field.

We start with the equation given in [24] for small radial departures of a plasma column from the position of equilibrium:

$$\frac{d^2\xi_r}{dt^2} = A\xi_r; \tag{4.1}$$

where

$$A = \frac{1}{4\pi a^2\rho}\left[H_\varphi^2 + \frac{(\varkappa a H_{ze} \pm m H_\varphi)^2 K_m(\varkappa a)}{\varkappa a\,(dK_m(\varkappa r)/d\varkappa r)_{r=a}} - \frac{\varkappa a H_{zi}^2 I_m(\varkappa a)}{(dI_m(\varkappa r)/d\varkappa r)_{r=a}}\right], \tag{4.1a}$$

a is the radius of the plasma column (the cross section of the column is assumed to be constant and circular independently of the method of stabilization), ρ is the plasma density, $\kappa = 2\pi/l$ is the wave number characterizing the kink deformation along the z axis with wavelength l. H_φ is the magnetic field on the undisturbed boundary of the plasma column due to the pinch current, H_{ze} is the longitudinal magnetic field outside the plasma column, and H_{zi} is the same field inside the plasma. In analyzing the perturbations, it was assumed that the plasma was perfectly conducting, and the perturbation of the plasma boundary was taken in the form $\xi_r = \xi(t)\exp(im\varphi + i\kappa z)$, where m is the mode number for the instability in the azimuth φ, $I_m(\kappa r)$ and $K_m(\kappa r)$ are m-th order Bessel functions of an imaginary argument of the first and second kind, respectively.

The stability condition for the system described by this equation reduces to a well-known criterion [24] which, for the most dangerous instability modes $m = 0$ (sausage) and $m = 1$ (kink and screw) reduces to

$$H_\varphi < \sqrt{\varkappa a\left[\frac{H_{ze}^2 K_0(\varkappa a)}{K_1(\varkappa a)} + \frac{H_{zi}^2 I_0(\varkappa a)}{I_1(\varkappa a)}\right]} \tag{4.2}$$

when $m = 0$, and to

$$H_\varphi < \frac{K_1(\varkappa a)}{K_0(\varkappa a)} \times\left[H_{ze} \pm \sqrt{H_{ze}^2 + \left\{H_{ze}^2 + \frac{I_1(\varkappa a)}{K_1(\varkappa a)}\,\frac{K_0(\varkappa a) + [K_1(\varkappa a)/\varkappa a]}{I_0(\varkappa a) - [I_1(\varkappa a)/\varkappa a]}H_{zi}^2\right\}\frac{\varkappa a K_0(\varkappa a)}{K_1(\varkappa a)}}\right] \tag{4.3}$$

when $m = 1$.

This condition becomes increasingly difficult to satisfy as the wavelength increases, i.e., κa decreases. In the limit as $\kappa a \ll 1$, this inequality reduces to

$$\frac{H_{ze}}{H_\varphi} > \frac{2}{\varkappa a\sqrt{1 + (H_{zi}^2/H_{ze}^2)}} \tag{4.4}$$

for the most unfavorable case when we take the negative sign in (4.3). For small κa this predicts a longitudinal field which is too high and may lead to fundamental difficulties even apart from technological difficulties.

For example, it is well known that, in high magnetic fields, magnetic radiation by electrons produces a substantial contribution to the energy balance.

The problem is then to reduce, one way or another, the stabilizing field. One possibility is to use high-frequency magnetic fields. There are a number of possibilities here, each of which must be considered separately. We shall follow the historical sequence of events and begin by considering the stabilization method illustrated in Fig. 9.

The Z-pinch with stationary current J is stabilized by high-frequency fields due to the currents in the multipole which is located on the periphery of the discharge chamber. The force per unit length of the pinch due to the φ component of the high-frequency field $\tilde{H}\cos\omega t$ in the r direction is $(J\tilde{H}\cos\omega t)/c$, or if we express the current in terms of the field H_φ due to this current on the boundary of the pinch, $(aH_\varphi\tilde{H}\cos\omega t)/c$. When the plasma column departs from the position of equilibrium by a small amount ξ_r in the radial direction, it experiences a restoring force $(a/2)\cos\omega t[H_\varphi(\partial\tilde{H}/\partial r)_a + H(\partial H_\varphi/\partial r)_a]\xi_r$. The second term in this expression is much less than the first, because the high-frequency field falls rapidly from the periphery toward the center, and relatively little on the boundary of the pinch. We shall therefore neglect the second term. There is, of course, no difficulty in taking this term into account but, for the sake of simplicity, we shall neglect it. We now substitute the remaining expression in (4.1) [we shall use the method mainly for the stabilization of kink-type perturbations and will investigate (4.1) for the case $n=1$]. The result of this is the Mathieu equation

$$\frac{d^2\xi_r}{dt^2}=\left[A+\frac{H_\varphi(\partial\tilde{H}/\partial r)_a\cos\omega t}{2\pi a\rho}\right]\xi_r. \tag{4.5}$$

When the constant component on the right-hand side is greater than zero, this is the equation for an inverted pendulum on an oscillating suspension. We shall only be interested in this particular case since the opposite situation means that (4.3) is satisfied, i.e., the system is stabilized by constant fields and there is no point in introducing high-frequency fields.

The stability criteria for the inverted pendulum were first obtained by P. L. Kapitza[25] (see also [26]) and by Jeffreys.[27] Their results are identical and the stability conditions for (4.5) reduce to

$$\omega^4>\left[\frac{H_\varphi(\partial\tilde{H}/\partial r)_a}{2\pi a\rho}\right]^2>2\omega^2 A. \tag{4.6}$$

From these two inequalities we obtain the following condition for the stabilizing field:

$$\left(\frac{\partial\tilde{H}}{\partial r}\right)_a>\frac{4\pi a\rho A}{H_\varphi} \tag{4.6a}$$

and for the frequency:

$$\omega>\sqrt{2A}. \tag{4.6b}$$

If stabilization is produced only by the high-frequency field due to the multipole, and there are no longitudinal stationary fields, then for $n=1$ condition (4.6a) assumes the form

$$\left(\frac{\partial\tilde{H}}{\partial r}\right)_a>\frac{H_\varphi}{a}\,\frac{K_0(\varkappa a)}{K_0(\varkappa a)+[K_1(\varkappa a)/\varkappa a]}, \tag{4.7}$$

which, for long waves, when $\kappa a\ll 1$ takes the form

$$\left(\frac{\partial\tilde{H}}{\partial r}\right)_a>\frac{H_\varphi}{a}\varkappa^2\ln\frac{2}{\varkappa a}=\frac{2J}{c}\left(\frac{2\pi}{l}\right)^2\ln\frac{l}{\pi a}. \tag{4.7a}$$

This form of stability condition is given in [22].

To ensure that the plasma column is stable for $m=1$ at all wavelengths, including $\kappa a\gg 1$, we have from (4.7)

$$\left(\frac{\partial\tilde{H}}{\partial r}\right)_a>\frac{H_\varphi}{a}=\frac{2J}{ca^2}. \tag{4.7b}$$

This condition can be satisfied only by very strong high-frequency fields, which means that there are major practical difficulties in using this method to stabilize appreciable currents. The second criterion, which defines the minimum frequency of the high-frequency field, is obtained from (4.6b) by assuming that $H_{zl}=H_{zi}=0$. Here,

$$\omega>\frac{H_\varphi}{\sqrt{4\pi a^2\rho}}\sqrt{\frac{K_0(\varkappa a)}{K_0(\varkappa a)+[K_1(\varkappa a)/\varkappa a]}},$$

or, since $H_\varphi=v_{Ti}(4\pi\rho)^{1/2}$ [where the thermal velocity of the ions is $v_{Ti}=(2T/M)^{1/2}$]

$$\omega>\frac{v_{Ti}}{a}\sqrt{\frac{K_0(\varkappa a)}{K_0(\varkappa a)+K_1(\varkappa a)/\varkappa a}}, \tag{4.8}$$

which for long waves yields

$$\omega>\frac{2\pi v_{Ti}}{l}\sqrt{\ln\frac{l}{\pi a}}, \tag{4.8a}$$

and for short waves

$$\omega>\frac{v_{Ti}}{a}. \tag{4.8b}$$

These inequalities clearly show that this is the most effective method for the stabilization of long-wave perturbations, whereas constant longitudinal fields stabilize short-wave perturbations quite easily.

At first sight, it might appear that the optimum system is that in which a constant longitudinal field stabilizes short-wave perturbations and a multipole high-frequency field suppresses long-wave perturbations. This is not so. It is shown in [28] that the combination of a longitudinal field and a high-frequency field will satisfy the stability conditions only if H_{ze} is close to the value for which (4.4) is satisfied, i.e., the stabilizing effect of the high-frequency field is then quite weak. This can be readily understood by considering the basic ideas involved in the stabilization and confinement of plasma by high-frequency fields. The starting point must be to avoid introducing a special direction in which the deformation may develop differently in comparison with other directions. The presence of an external stationary field leads to the appearance of a direction which depends on the sign in front of m, i.e., on the direction in which the pinch is twisted. The stability conditions are different, depending on whether this twisting action is left- or right-handed relative to the field and, since the effect of the hf field is, on the average, independent of the direction of twisting, this field turns out to be ineffective. The field inside the plasma does not interact with the external high-frequency field and, therefore they facilitate the improvement in the stability conditions.

In fact, when $H_{ze}=0$ and $H_{zi}\neq 0$, condition (4.6a) yields

$$\left(\frac{\partial\tilde{H}}{\partial r}\right)_a>\frac{H_\varphi}{a}\,\frac{K_0(\varkappa a)}{K_0(\varkappa a)+[K_1(\varkappa a)/\varkappa a]}\times\left\{1+\frac{H_{zi}^2}{H_\varphi^2}\,\frac{I_1(\varkappa a)}{K_1(\varkappa a)}\,\frac{K_0(\varkappa a)+[K_1(\varkappa a)/\varkappa a]}{I_0(\varkappa a)-[I_1(\varkappa a)/\varkappa a]}\right\}. \tag{4.9}$$

To each value of H_{zi}^2/H_φ^2 there corresponds a certain κa for which the right-hand side of this inequality is a maximum. The dependence of this maximum on H_{zi}^2/H_φ^2 is shown in Fig. 10, which also shows some values of κa corresponding to the maximum. For example, it follows from the figure that when $H_{zi}^2/H_\varphi^2=0.5$ the value of $(\partial\tilde{H}/\partial r)_a a/H_\varphi$ for which stability persists for the $m=1$ mode amounts to about 0.1 for all wavelengths, and it follows from (4.2) that for this field ratio all wavelengths with $m=0$ are also stable.

These calculations suggest that a system incor-

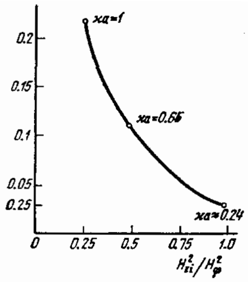


FIG. 10. Dependence of the right-hand side of (4.9) on H^2_{zi}/H^2_φ, for values of κa for which, with given H^2_{zi}/H^2_φ, this function assumes the maximum value.

porating a longitudinal field localized inside the pinch and an hf field due to a multipole has definite advantages.

It was pointed out above that the stability conditions depend on the direction in which the pinch is twisted because (4.3) includes two solutions, each corresponding to a particular direction of twist relative to the external field. In one case, (4.3) reduces for long waves to the well-known criterion

$$\frac{H_{ze}}{H_\varphi} > \frac{2}{\varkappa a\sqrt{1+(H^2_{zi}/H^2_\varphi)}}, \tag{4.10a}$$

and, in the other case, it reduces to

$$\frac{H_{ze}}{H_\varphi} > \frac{\varkappa a}{2}\ln\frac{2}{\varkappa a}. \tag{4.10b}$$

It is clear that the second condition is much more readily satisfied than the first.

When either the external field or the current in the pinch are of high frequency, it is natural to expect that the stability conditions will correspond to an intermediate position between the two limiting values defined by these two inequalities. This system with a high-frequency current in the pinch and constant longitudinal fields was proposed and investigated in [29]. The stability condition for this case can be written in the form

$$\left(\frac{H_{ze}}{\widetilde{H}_\varphi}\right)^2 > \frac{K_0(\varkappa a)}{2\varkappa a K_1(\varkappa a)}\left\{1+\left(\frac{H_{zi}}{H_{ze}}\right)^2\frac{I_1(\varkappa a)}{K_1(\varkappa a)}\,\frac{K_0(\varkappa a)+[K_1(\varkappa a)/\varkappa a]}{I_0(\varkappa a)-[I_1(\varkappa a)/\varkappa a]}\right\}^{-1}, \tag{4.11}$$

which for long waves with $\kappa a \ll 1$ becomes

$$\left(\frac{H_{ze}}{\widetilde{H}_\varphi}\right)^2 > \frac{\ln(2/\varkappa a)}{1+(H^2_{zi}/H^2_{ze})}. \tag{4.11a}$$

The frequency condition for this system is

$$\omega > \frac{\widetilde{H}_\varphi}{\sqrt{4\pi a^2\rho}}; \tag{4.12}$$

where $\widetilde{H}_\varphi$ is the field at the edge of the pinch due to the high-frequency current flowing through the discharge.

It is readily seen that condition (4.11a) is the square of the geometric mean of conditions (4.10a) and (4.10b) for the effective high-frequency field. It is clear that since $\widetilde{H}_\varphi$ is produced by the high-frequency current, the special direction is eliminated because twisting in either direction is equally probable. Therefore, the effect averaged over both directions is operative in this case.

There is, of course, very little change in the situation when the field H_z is a high-frequency field and the current is constant. The critical conditions remain the same as before but the possibility of resonances must be borne in mind in both these variants.

It is well known that in systems described by equations with periodic coefficients, the regions of stability alternate with regions in which the system is unstable. Particularly dangerous are those conditions which lead to the appearance of parametric resonances occurring when the natural frequency of oscillation of the plasma column approaches half the frequency of the hf field. These questions are considered in [30].

The present author has investigated experimentally dynamic stabilization by an hf longitudinal field (see [31,32]). Unfortunately, these studies have not been continued. Insofar as the appearance of parametric resonances in such a system are concerned, the experiments reported in [33,34] did reveal the presence of a phenomenon which could be treated as parametric excitation. At any rate, oscillations at half the frequency were observed. However, this does not mean that there is little point in pursuing investigations in this direction because it would appear to be quite easy to provide conditions under which parametric resonances can be avoided.

As an example, let us return to the experiments with a traveling wave. In this system, there is a longitudinal drift current in the high-frequency field of the traveling wave, so that the conditions obtaining in this case are close to the ideas developed in [29]. No evidence has been obtained for the presence of instabilities, at any rate the MHD type, and, even more so, parametric resonances.

We must now return to (4.11a). For this case, $H_{zi} = 0$, $\kappa_{min} = 2\pi/l_{max}$, so that

$$\varkappa_{min}a = \frac{a}{R} \approx 0.1,\quad \left(\frac{\widetilde{H}_{ze}}{H_\varphi}\right)^2 > \frac{\ln(2/\varkappa a)}{2} \approx 1.5.$$

According to the results given above, $(H_\varphi/\widetilde{H}_z)_{max} \approx 0.8$ and, consequently $(\widetilde{H}_{zl}/H_\varphi)^2_{min} \approx 1.55$ which satisfies the above criterion. This suggests that the stability of the drift current in the traveling-wave system is ensured by the stabilizing effect on the hf field. Moreover, this enables us to conclude that the relationships obtained for systems of this kind are at least to some extent confirmed by the results of the above experiments.

Theoretical studies of various aspects of the above variants of dynamic stabilization are considered in [25-40]. Our references do not pretend to be exhaustive and include only some of the published papers.

5. DYNAMIC STABILIZATION: RESULTS OF EXPERIMENTS AND CONSEQUENCES

We shall describe experiments on the stabilization of the Z-pinch by high-frequency multipole fields. In these experiments, longitudinal fields were not employed. The high-frequency fields were produced by a multipole located on the periphery of the discharge chamber in such a way that the currents in neighboring rods forming the multipole were shifted in phase by 180°. Figure 11 illustrates the hexapole arrangement. Two series of experiments were performed.

In the first series, the multipole rods were supplied by a damped circuit consisting of capacitors with relatively high Q ($Q \approx 60-80$).

In these experiments we used six stabilizing rods. The discharge chamber was in the form of a glass cylinder, 20 cm in diameter and 60 cm long, covered by flat electrodes at each end. Gas was exhausted and admitted to the system through one of the electrodes. The stabilizing circuit was in the form of a set of six

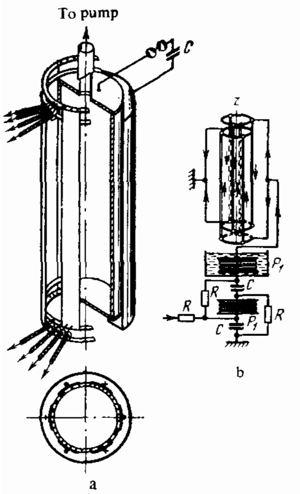


FIG. 11. Schematic diagram of the discharge chamber (a) and supply circuit for dynamic stabilization (hexapole) in a system fed by a damped circuit (b).

bus bars arranged, as shown in the figure, on the outer surface of the chamber. This system was connected by 36 cables to a bank of high-quality capacitors with a total capacitance of 0.6 μF. The connections were made so that at each instant of time the currents flowing through neighboring bus bars were opposite in direction. At a frequency of about 700 kHz the amplitude of the current on each bus bar was about 27 kA.

The Q factor of the stabilizing circuit with the discharger and plasma was $Q \approx 20$. The parameters of the main discharge-current circuit were $C = 70\ \mu F$, $U = 8-12$ kV. The period of the discharge circuit was about 40 μsec.

The main diagnostic means employed in this cycle of experiments[14,41-43] were the following: streak photography of the discharge, measurements of the longitudinal field H_z which appeared as a result of loss of stability with $m = 1$ (kink and screw instabilities), and measurements of the energy flux to the chamber walls.

The photographic method was used to record the luminous diameter of the plasma column and provided an estimate of the time of its localization near the chamber axis. The slit of the photographic recorder was perpendicular to the axis of the discharge chamber. One of the bus bars contained longitudinal gaps and the observations were carried out through these gaps and the spaces between the bus bars. The field component H_z was measured with a single-turn coil around the chamber. The integrated signal was applied to oscillograph plates so that the signal amplitude was proportional to the component of the discharge current producing the field H_z.

The thermal probe described in [44] was used to record the beginning of the interaction between the plasma and the wall and to determine the rate of escape of energy from the plasma column. This method enabled a reliable recording to be made of the energy flux up to 5×10^{-3} J/cm^2.

In all the measurements, an oscillograph was used to monitor the phase at which the stabilizing voltage was switched on. In this way, it was possible to achieve simultaneous synchronous recording of all the quantities listed above. The photographic time scan of the emission (Fig. 12), taken from,[42] shows that, when both the main and the stabilizing circuits are acting

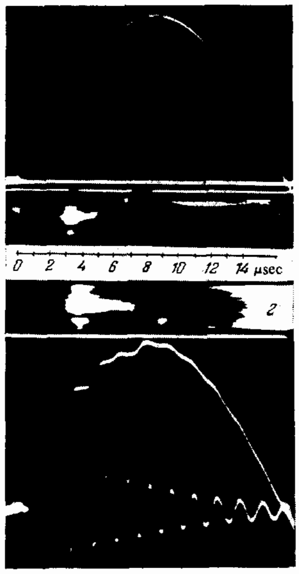


FIG. 12

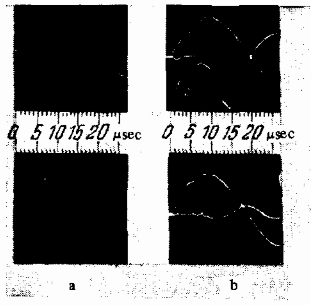


FIG. 13

FIG. 12. Oscillograms of the discharge current and the stabilization current, and photographic time scan of the emission from the plasma column in the absence (1) and presence (2) of stabilization, obtained for a system with a damped circuit.

FIG. 13. Oscillograms of thermal probe signal (a) and longitudinal field signal (b). Upper trace - without stabilization; lower trace - with stabilization. These oscillograms correspond to the photographs and oscillograms in Fig. 12 (1-2).

simultaneously, the column is confined near the axis for a time interval of about 6 μsec, whilst, in the absence of stabilization this time is much shorter.

Figure 13a shows the thermal-probe signal synchronized with the discharge current, and Fig. 13b shows the oscillograms of the field H_z. It is clear from the photographs in Fig. 12 that the instant of appearance of the field H_z corresponds to the disintegration of the column. It is precisely at this moment that the emission localized near the axis is found to disappear. It is important to note that, in the absence of stabilization, the plasma column will disintegrate under these conditions almost immediately after it is formed or, more precisely, at the time when the characteristic second break is observed on the current curve (this is the second singularity; see, for example, [45]).

When the stabilizing field is imposed, the time of disintegration of the column is shifted so that disintegration is observed on the maximum and sometimes even after the maximum of the discharge current, when the stabilizing field is already subject to considerable damping. The thermal-probe signals (Fig. 13) which are synchronized with the discharge-current oscillogram clearly show the delay in the arrival of the energy flux on the wall in the presence of stabilization, as compared with the unstabilized discharge.

It may be concluded that the onset of intensive arrival of the energy flux on the wall corresponds to the instant of disintegration established by other methods, i.e., by photographs of the emission and measurements of H_z.

Comparison of the results of all these measurements leads to the conclusion that, for certain definite values of the parameters characteristic for devices of this kind, there is in fact a well-defined stabilization effect.

In a relatively broad range of gas (hydrogen) pressures, roughly between 0.05 and 1 Torr, the plasma column is confined to the axis of the chamber, with a mean diameter of 3 to 3.5 cm, up to the instant of time when the current reaches about 100–150 kA. At the same time, the stabilizing current is reduced by a factor of about two owing to damping.

During this time the column does not kink (there is no H_Z) and the thermal probe does not show the presence of any appreciable arrival of energy at the wall. It is important to note that stabilization is not observed if the hf circuit is switched on after the beginning of the contraction process, when the column has already begun to deform. This method of stabilization merely prevents the appearance of instabilities but cannot suppress an already existing instability. This situation was observed in all the experiments.

In order to obtain a clearer picture of the development of instabilities and of the stabilization effect, we carried out experiments on installations with somewhat different parameters but essentially similar to those described above. In these experiments, the plasma column was examined by the framing camera. In this way, one could obtain photographs of the entire plasma column at intervals of about 1 μsec and thus establish the deformation of the column, the wavelength of the instability, and the deformation of the cross section of the column. In the first series (see [41]), photographs recorded on the side wall were obtained (see Fig. 14).

It is clear from these photographs that, in the absence of stabilization, the column breaks up as a result of kink and screw instabilities, and the deformation begins with long-wave perturbations which eventually develop into shortwave perturbations. When the stabilizing circuit is switched on, there are no visible instabilities when the circuit parameters satisfy certain definite conditions.

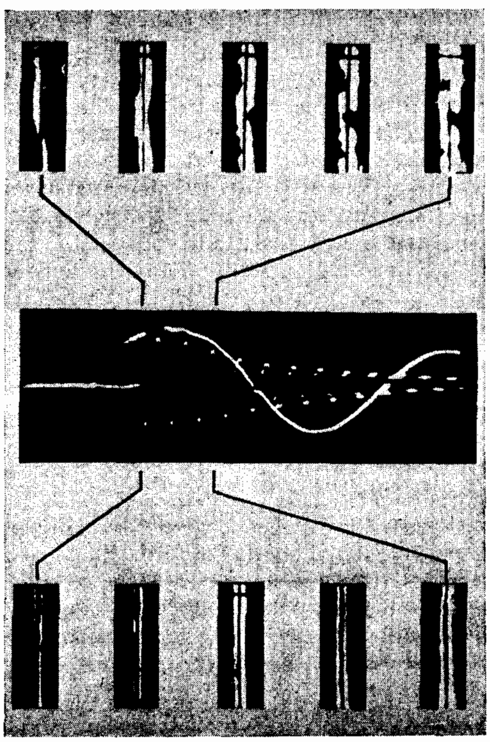


FIG. 14. High-speed photographs of the Z-pinch recorded through the side wall (framing camera; 1 frame per μsec) and synchronized with the discharge and stabilizing currents. Top - photographs of the plasma column in the absence of stabilization; bottom - ditto in the presence of stabilization. The interval of time during which the photographs were recorded is indicated. The shadow at the center of column is the projection of the stabilizing rods.

The lifetime of the stable state does, of course, depend on the damping of the current in the stabilizing circuit. This time is found to increase with increasing Q of the circuit, as shown in [41].

The results of experiments with a damped stabilizing circuit given here are subject to an important shortcoming.

The stabilization is carried out mainly in the region of growing current, while the amplitude of the stabilizing current decreases. The result of this is that, as the discharge current increases, there is a reduction in the stabilizing field. It is therefore natural to try and investigate dynamic stabilization with an undamped circuit supplied by a vacuum-tube generator. Although the technological facilities at our disposal place a substantial restriction on the magnitude of the stabilizing current, nevertheless, we have succeeded in obtaining results of definite interest.

The experiments (see [46–48]) were carried out on an installation consisting of a glass discharge tube, 6.5 cm in diameter, in which the distance between the electrodes was 75 cm. One of the electrodes had an axial aperture through which observations could be carried out. The electrodes were connected to a capacitor bank of 22 μF and the anode voltage applied to them could be varied between 8 and 22 keV. The switching was carried out through an air spark gap. The inductance of the circuit was about 8 mH. Under these conditions, the discharge current could be varied from 12 to 35 kA, and the initial rate of rise of the current $(dJ/dt)_0$ from 1.5 to 4.5×10^9 A/sec. One quarter of the period, i.e., the time taken for the current to rise to its maximum, was 19 μsec.

A quadrupole was used for stabilization. The length of the rods was 70 cm and their diameter 1 cm. The distance between opposite rods was 10 cm. The oscillatory circuit consisted of the stabilizing rods with an equivalent inductance of 0.39 μH and a bank of ceramic capacitors of 54×10^{-3} μF, so that the working frequency was about 1 MHz. The effective Q of the circuit in the course of the discharge was $Q\approx25$. The circuit was supplied through a matching unit from a pushpull self-excited oscillator consisting of four vacuum tubes. For a steady output power of about 4 MW, about 2 MW was transferred to the circuit. The current in each rod could then be increased up to 3.5 kA. The corresponding radial rate of increase in the field was $\partial H/\partial r\approx110$ Oe/cm.

For diagnostic purposes, we used both framing and streak cameras, magnetic probes for measuring the high-frequency fields and the field due to the discharge current, and coils for the detection of the longitudinal magnetic field H_Z. Moreover, from the position of the limit of the continuous spectrum in the visible part of hydrogen emission, it was possible to estimate the mean charged-particle concentration.

Main results. The plasma column is formed only after a certain definite rate of increase of the current has been reached. In the above experiments, this minimum rate of increase was $\partial J/\partial t\geq2\times10^9$ A/sec. Moreover, the formation of the column was very dependent on the delay time Δt, i.e., the interval between the moment when the high-frequency stabilizing field was switched on and the moment at which the main discharge was switched on. Stable confinement of the plasma column was observed for $40 < \Delta t < 100$ μsec, and all the results corresponded to this range of delay time. For $\Delta t < 40$ μsec, the plasma shrinks to the axis too rapidly, and for $\Delta t > 100$ μsec the plasma column is not formed. This is, of course, connected with the degree of ionization of the gas at the beginning of current flow. When the degree of ionization is low, the process occurs too rapidly and, as will be shown below, short-wave instabilities develop and cannot be stabilized by the methods employed here. For $\Delta t > 100$ μsec, this current is insufficient for compressing the gas because of the presence of heavy gas impurities due to the evaporation of the wall material. The experiments were carried out for an initial gas (hydrogen) pressure of $p_0 = 0.2$ Torr. This pressure corresponds to the most favorable conditions for the processes in which we are interested. Unfortunately, the most important parameter is the initial rate of increase in the discharge current which must not lie outside the relatively narrow range $2 \times 10^9 < (\partial J/\partial t)_0 < 3.5 \times 10^9$ A/sec. If the upper limit is exceeded, the plasma column disintegrates very rapidly for reasons which will be indicated below. If the lower limit is not reached, the plasma does not separate from the walls.

When all these conditions are satisfied, a plasma column is clearly formed (Fig. 15) and is stabilized on the axis for a certain interval of time. For comparison, Fig. 16 shows photographs of the discharge obtained under the same conditions but without the stabilizing field.

We reproduce streak and framing camera photographs recorded on the side wall and from the end of the chamber, and oscillograms of the discharge current and longitudinal magnetic field H_z generated during the development of the instability. The electron density determined from the Balmer series limit is not less than 10^{16} cm^{-1}. The radius of the compressed column determined from the framing camera photographs is ~0.5 cm under these conditions. The lifetime of the stable compressed state does not exceed ~10 μsec under these conditions.

In order to explain the limited lifetime of the pinched state, it is not necessary to assume that the plasma column became unstable after a certain interval of time. As in the case of the traveling wave, this time is restricted by the overheating process.

We shall now use the well-known formula $J = 4c^2NT$ (valid for the Z-pinch) which is deduced from the balance of the magnetic-field and gas-kinetic pressures. In our case of relatively low currents, when the temperature does not exceed ~10 eV and $J = 20$ kA (this is close to the maximum value for which the breakup of the plasma column is observed), we find that $N = 6 \times 10^{16}$ cm^{-1}.

The total number of particles per unit length of the discharge tube at a pressure of 0.2 Torr was $\sim 5 \times 10^{17}$ cm^{-1}.

Thus, as the particles are captured into the plasma column, become ionized, and heated, the plasma pressure becomes such that the current cannot confine these particles by its field, and the column breaks up. According to the above estimates, the time of ionization and

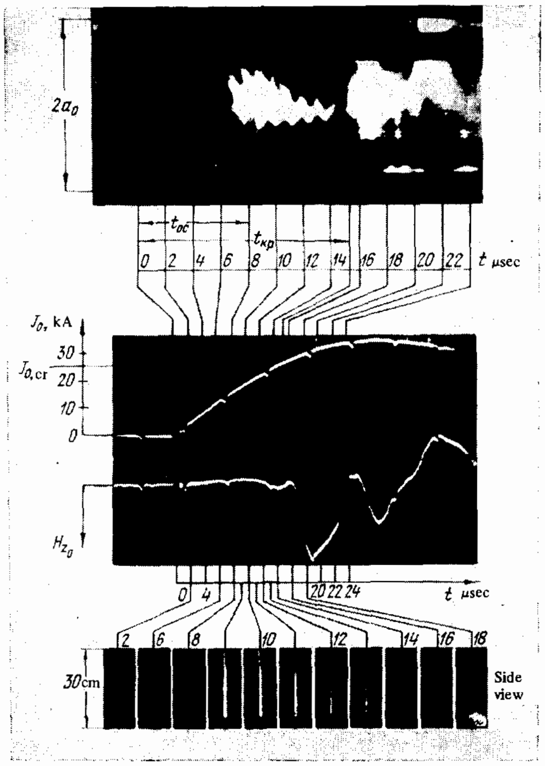


FIG. 15. Experiments with undamped stabilizing circuit fed by a vacuum-tube generator. Top - streak photograph of the plasma column. Bottom - framing camera photographs (1 frame per μsec) taken from the side wall. The middle photograph shows oscillograms of the discharge current J and longitudinal field H_z. The quantity J_{crit} is the current for which the plasma column disintegrates; a_0 is the internal radius of the discharge tube ($\tilde{J} = 2.8$ kA, $p_0 = 0.2$ torr. $J_{0_m} = 35$ kA, $J_0(t = 0) = 2.4 \times 10^9$ A/sec).

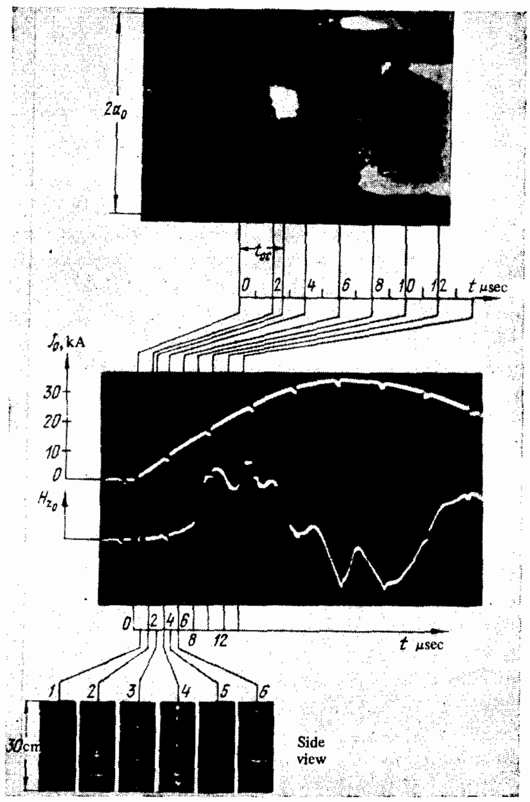


FIG. 16. As in Fig. 15 but without stabilization (D_2, $p_0 = 0.2$ Torr, $J_{0_m} = 35$ kA, $J_0(t = 0) = 2.4 \times 10^9$ A/sec).

heating of the gas entering the column from the periphery is quite close to the observed lifetime of the stable state.

The question why the lifetime of the stable state is so short can therefore be answered relatively simply. On the other hand, it is much more difficult to explain why this state exists at all. For example, are there no instabilities with $m = 0$, and short-wave instabilities with $m = 1$, for which the stability criteria introduced above are definitely not satisfied?

Let us now consider the minimum wavelength for which the stability conditions are satisfied. In the above experiments $(\partial H/\partial r)_0 = 100$ Oe/cm, $J = 2\times 10^4$ A, and $a = 0.5$ cm. We shall see that, in this case, we can use (4.9a) from which it follows that $l_{min} \geq 80$ cm, i.e., the kink and screw deformations are stabilized for wavelengths close to the separation between the electrodes. A roughly similar situation occurs in experiments with the damped circuit, described above, and in other investigations of this kind.[14,42,43,45,49]

We have thus encountered a problem which is relatively rare in practical physics. A more common situation is that in which one has to explain the reasons why the experimental results disagree with the theoretical predictions to an extent greater than one would expect in view of the adopted assumptions. The theory is usually better than the experiment. Here, we have the opposite situation. The results of the experiment are more satisfactory than is predicted by the theory. It turns out that it is sufficient to satisfy the stability conditions from some relatively long wavelength onward and, when this is so, short-wave perturbations and $m = 0$ perturbations are not observed.

From the standpoint of the theory of magnetohydrodynamic stability, developed for stationary plasma and used above, this is a contradictory situation. Instabilities of this kind develop more readily than shorter wavelengths. The growth rate is known to be inversely proportional to the wavelength. One would therefore expect to observe the appearance and development of short-wave instabilities with $m = 1$ and $m = 0$ modes.

The above results clearly indicate that we are dealing with phenomena which do not fit into the framework of the results obtained for stationary Z-pinches. It is therefore natural to ask: what governs the instability wavelength in the unstabilized Z-pinch? Which discharge parameters determine this quantity? To obtain an answer to these questions, we have carried out experiments in which the development of instabilities could be observed in a broad range of values of the parameters characterizing the discharge. In these experiments, we varied the geometric dimensions (length of discharge gap L and tube radius a_0), the initial rate of increase in the current $(dJ/dt)_0$, and the initial gas (hydrogen) pressure p_0. In all other respects, the installation was the same as that described above. The wavelength was determined with the aid of a high-speed framing camera.

Figure 17 shows (top row) photographs of the column for discharge tubes of different length. The chosen frame corresponds to the instant of time close to the maximum compression of the column. It is clear that there is no substantial difference between the wavelengths. Hence, it can be concluded that the length of the discharge gap has no appreciable effect on the instability wavelength.

A completely different situation is observed in experiments in which the values of $(dJ/dt)_0$ and the radius

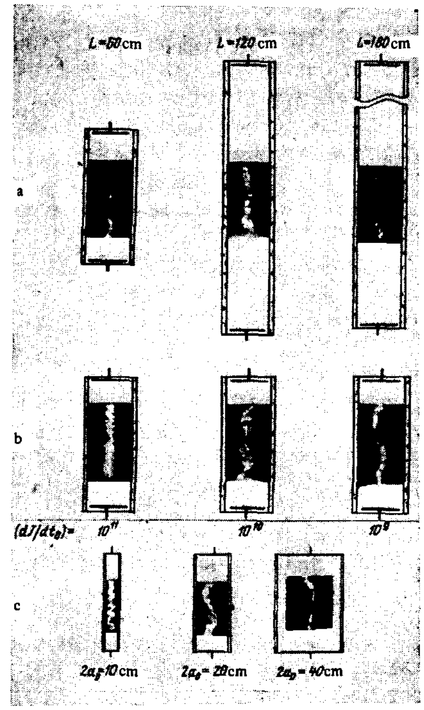


FIG. 17. Typical photographs obtained with a framing camera (p_0 = 0.1 mm Hg of H_2): (a) dependence of wavelength on the length of the discharge chamber [$2a_0$ = 20 cm, $(dJ/dt)_0 = 10^9$ A/sec]; (b) ditto as a function of the rate of increase in the current (h = 60 cm. $2a_0$ = 20 cm); (c) ditto as a function of the initial radius [l = 60 cm, $(dJ/dt)_0 = 5\times 10^9$ A/sec].

a_0 of the discharge tube varied. The second row in Fig. 17 shows the corresponding photographs. It is clear that the instability wavelength increases with increasing a_0 and $(dJ/dt)_0$. The experiment was performed for $10^9 < (dJ/dt)_0 < 10^{11}$ A/sec. For $(dJ/dt)_0 > 10^{11}$ A/sec the $m = 1$ instabilities go over into $m = 0$ instabilities.

Using the results of experiments performed for a relatively broad range of values of $(dJ/dt)_0$ and a_0, it is possible to establish a definite relationship between the instability wavelength and the values of $(dJ/dt)_0$ and a_0. It turns out that, in the range $5\times 10^9 < (dJ/dt)_0 < 10^{11}$ A/sec and for a_0 between 4.5 cm and 20 cm, the product $l(dJ/dt)_0^{1/2}/a_0$ varies very little and remains approximately equal to 2×10^5. Hence $l \approx 2\times 10^5 a_0 (dJ/dt)_0^{-1/2}$.

We note that no clearly defined effect was observed due to variations in the initial pressure. It would appear that the dependence on this parameter is very slight. Most of the measurements were performed for $p_0 \approx 0.1$ Torr. The onset of instability was fixed by two magnetic probes which measured the longitudinal field component H_Z. One probe was located near the walls (this is the H_Z probe) and the other (H_{Z_0}) was placed on the axis of the chamber. It was found that the wall probe was the first to record the appearance of the instability Hence, it follows that the instability sets in at a relatively early stage of the compression process, quite close to the instant of separation of the plasma column from the chamber walls. This is confirmed by the fact that the dynamic stabilization operates effectively only if the stabilizing pulse is applied before the column contracts toward the axis.

The results given above cast doubt on the validity of the basic assumptions on the nature of instabilities appearing in relatively fast Z-pinches. A clear dependence of the instability wavelength on the initial rate of increase in current and tube radius are not consistent with the description involving MHD-type instabilities in the stationary Z-pinch for which the current and radius of the compressed plasma column are characteristic quantities. Here, as we have seen, the instability is connected with the dynamics of the column compression process and not with its appearance in the compressed state.

We shall now attempt to explain the observed contradiction as follows. It is pointed out above that, in our case, the dynamics of the process plays the main role and there is hardly any point in considering stability as the development of a perturbation of some quasistationary equilibrium state. A more natural assumption is that the instability is of the Rayleigh-Taylor type and arises as a result of the acceleration of the plasma during its motion from the periphery toward the center. The instability may develop prior to the separation of the entire plasma from the walls of the discharge tube, as a result of the propagation of a compression wave. There is a long-standing proposal that, in fast Z-pinches, the $m = 0$ instabilities are of the Rayleigh-Taylor origin. Experiments have been performed which have provided relatively satisfactory confirmation of this hypothesis. Theoretical studies of such processes, among which we note the work of Harris,[50] show that this instability will also occur for the $m = 1$ mode. Harris has also given a bibliography of experimental studies.

However, the assumptions upon which Harris's work is based have led to the conclusion that perturbations of all wavelengths are stable. This result is explained by the adopted model.

We shall summarize here the results reported in [51] and based on an analysis of certain dimensional relationships. The starting point was the well-known Kurskal and Schwarzschild relation obtained from an analysis of Rayleigh-Taylor instabilities in semi-infinite plasma confined by a magnetic field in a gravitational field:

$$\omega = \frac{1}{2}\sqrt{g\varkappa}, \tag{5.1}$$

where g is the acceleration of the boundary of the plasma, κ is the wave number of the perturbation harmonic, and ω is the instability growth rate.

Since the above experimental results establish the existence of a clear dependence of the instability wavelength on the radius of the discharge tube, it is natural to assume that this instability develops while the plasma still interacts with the walls of the discharge chamber, i.e., before the plasma contracts toward the axis. The only quantity which determines the time scale is the time of separation of the plasma from the walls, t_{sep}. We shall suppose that $\omega = 1/t_{sep}$. This is a purely dimensional relationship which we introduce without sufficient justification.

The acceleration of the plasma boundary during the initial stage of the compression process can be determined from the equation of motion for a mass of gas during compression by the magnetic field due to the current flowing through the plasma (see, for example, [49,52]):

$$g = \frac{(dJ/dt)_0}{a_0 c\sqrt{3\pi\rho_0}}. \tag{5.2}$$

During this stage, the current is assumed to be a linear function of time, i.e., $J = (dJ/dt)_0 t$. From (5.1) and (5.2), we can then find the instability wavelength by reducing all the quantities to the separation time:

$$l = \frac{J_{sep}^2}{2c\sqrt{MN}\,(dJ/dt)_0};$$

where M is the ion mass, N is the number of particles per unit length of the discharge chamber, $MN = \pi a_0^2\rho_0$, and J_{sep} is the current at the time of separation of the plasma from the vessel walls. This current was calculated in [45] and is given by

$$J_{sep} = c\sqrt[4]{\frac{80a_0^2\varepsilon\pi}{\sigma_{in}c}\left(\frac{dJ}{dt}\right)_0\sqrt{MN}}.$$

In [45] the separation current was determined for weakly ionized plasma. The ionization was produced during the development of the discharge, so that in the expression for the separation current we have the electron factor ϵ which we assumed to be approximately 100 eV per ionization event. The capture of particles while the plasma moves toward the axis of the chamber takes place for low degrees of ionization through the charge-transfer mechanism which of all the elementary events in this energy range has the largest cross section (in hydrogen $\sigma_{in} \approx 3\times10^{-15}$ cm^2).

If we now introduce the separation current into the expression for l, and substitute the various numerical values, we obtain in practicel units

$$l \approx \frac{a_0}{\sqrt{(dJ/dt)_0}}\,\frac{5\cdot10^9}{\sqrt[4]{N}}. \tag{5.3}$$

In the experiments reported in [51] there were $N \approx 10^{18}$ particles/cm, in which case $l \approx a_0(dJ/dt)_0^{-1/2}\cdot 2\times10^5$ which agrees with experimental results rather too well.

It is clear that, since the fourth root of N is present in (5.3), the dependence on the initial pressure has not been noticed, especially since the experiments were carried out in a relatively narrow range of initial pressure values.

It follows from the above results that the $m = 0$ perturbations may appear when $l/a_0 < 1$, and this means that

$$\left(\frac{dJ}{dt}\right)_0 > \frac{3\cdot10^{19}}{\sqrt{N}} \tag{5.4}$$

When $N = 10^{18}$, neck formation occurs for $(dJ/dt)_0 > 3\times10^{10}$, which is also in good (even too good) agreement with all the experimental results.

We shall now compare the above results with the experimental results on dynamic stabilization, assuming that to suppress instabilities it is sufficient to satisfy (4.9) at the time of separation of the plasma, i.e., when $a = a_0$, $J = J_{sep}$ and l is given by (5.3).

As an example, consider the experiment with the undamped circuit described above. For this case,

$$l = \frac{a_0}{\sqrt{\left(\frac{dJ}{dt}\right)_0}}\,\frac{5\cdot10^9}{\sqrt[4]{N}} \approx 16\text{ cm}, \varkappa a_0 = \frac{2\pi a_0}{l} = 1.25,$$

since $a_0 = 3.2$ cm.

Condition (4.9) assumes the form

$$\left(\frac{\partial\widetilde{H}}{\partial r}\right)_0 > \frac{0.2J_{sep}}{(3.2)^2}\,\frac{K_0(1.25)}{K_0(1.25)+[K_1(1.25)/1.25]} \approx 10^{-2}J_{sep},$$

where the current is in amperes.

The separation current in these experiments is ~10 kA and, as indicated above, the value of $(\partial\tilde{H}/\partial r)_0$ for which a well-defined stabilization effect is observed ~100 Oe/cm. It is clear that there is an equally surprising agreement between the above relationships and experimental data.

Analogous calculations performed for the experiments reported in [14,42,43,44] have led to similar results.

Such precise quantitative agreement between the results must be viewed with considerable scepticism, especially since the foundations upon which the calculations are based are decidedly shaky. Nevertheless, it is quite clear that these results force us to reconsider the qualitative explanation of the phenomena. Even when a theory becomes available which will provide an explanation for the appearance of this instability, and the mechanism responsible for the appearance of the instability during the early stage of compression of the pinch will be understood, it will still be unclear why the usual types of instability do not develop in the compressed pinch. The lifetime of the compressed state observed in all the experiments is much greater than the development time for MHD-type instabilities. The theory of these instabilities shows that the instability development time is proportional to the wavelength. Therefore, short-wave perturbations should develop in very short intervals of time ($<10^{-7}$ sec), but no appreciable deformations in the plasma column develop for at least several microseconds, which is much longer than the instability development time, even for the longest wavelengths. Nor can this anomalous instability be explained in terms of fields frozen into the plasma. It would appear that the high-frequency fields penetrate the plasma during an early stage of discharge development and then, as the process continues into the compression stage, the captured field somehow increases and somehow stabilizes the column. The most optimistic estimates of the attenuation time for the captured field show that the fields are damped out in a time which is much shorter than the lifetime of the compressed plasma column. None of these explanations is therefore satisfactory.

An explanation can only be obtained from a deeper understanding of the various phenomena, and this will only be possible after additional experimental studies.

The above results and speculations lead to the conclusion that studies of the above processes have, at least so far, led to more questions than answers.

Studies of the confinement of plasma by high-frequency fields and experiments on dynamic stabilization have produced such surprises that very little remains of the assumptions upon which the experiments themselves were formulated.

This, in turn, means that the many assumptions and their theoretical justification must be approached with considerable caution, at least until really satisfactory experimental data confirming or contradicting the initial assumptions become available.

There is, therefore, little point in considering in detail all the various existing suggestions. We note only those upon which we can pass a meaningful judgment.

6. HYBRID SYSTEMS

Hybrid (or composite) systems are defined as those in which the hf fields are not the main fields, not even for the suppression of MHD-type instabilities, but act as auxiliary fields in combination with constants or quasiconstant fields. Among the large number of proposals in this area, we consider only two.

We start with a brief description of the composite systems in the form of toroidal traps, known as RO and RT-2.[53–55] The differences between these two devices are not very substantial, at least insofar as the principles on which they are based are concerned.

These traps employ a strong quasiconstant magnetic field (10–20 kOe), and a high-frequency magnetic field is used to compensate the toroidal inhomogeneity and to heat and stabilize the hot plasma instabilities in the trap.

It is shown in [53] that it is possible to stabilize both MHD and kinetic plasma instabilities with rotating high-frequency fields, where the rotation of the field is carried out about the longitudinal axis of the torus and the field strength increases from the center toward the periphery.

Such fields can be produced by the helical multipole illustrated in Fig. 18, where the phase difference between neighboring conductors is $\pi/2$. However, the application of only the high-frequency field, even a helical field, would not ensure plasma equilibrium. Stellarator-type coils with quasiconstant currents are therefore used to compensate toroidal effects. The main parameters of the RT-2 installations are as follows.

The toroidal chamber has a major diameter of 130 cm and a minor diameter of 2.5 cm. The longitudinal magnetic field is 20 kOe and the pulse length is 0.01 sec. A double helical coil producing the quasiconstant field is employed with six steps along the major radius; the current in the coil is 100 kA.

The high-frequency field is produced by a self-excited vacuum-tube oscillator with a power output of about 60 MW at 2 MHz, which produces a pulse length of 0.001 sec. The hf field amplitude on the periphery is about 200 Oe. An octuple three-step winding is employed.

The installation is designed for experiments on plasmas with density $n \approx 10^{12}-10^{15}$ cm^{-3} and $T_e = 10-100$ eV. Under these conditions, $\bar{\beta} = 8\pi nT/\bar{H}^2 \ll 1$ and $\tilde{\beta} = 16\pi nT/\tilde{H}^2 \gtrsim 1$.

The first step was, of course, to determine the depth of penetration of the high-frequency field into the plasma. As usual, magnetic probes were used for this purpose. An anomalous penetration depth was observed which, again, cannot be explained in terms of existing theories. It was found that an increase in the quasiconstant longitudinal field H_z was accompanied by an increase

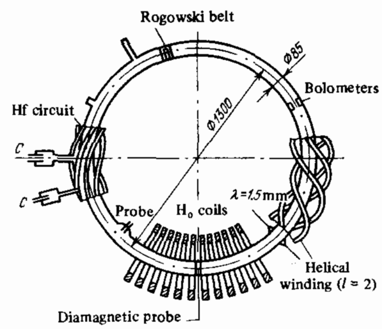


FIG. 18. Disposition of coils producing quasiconstant and hf fields in the RT-2 installation.

in the hf field at the center of the chamber. The same effect was observed when the initial gas pressure was reduced.

The high-frequency current in the plasma was practically independent of the type of gas employed (the measurements were carried out with hydrogen, helium and xenon in the chamber). It was also independent of the initial pressure and amounted to ~2.2 kA, which corresponded to a magnetic field of about 150 Oe on the boundary of the plasma column. Up to pressures $p_0 \approx 2\times 10^{-3}$ Torr, the skin layer was nearly classical. For $p_0 < 2\times 10^{-3}$ Torr the skin layer was anomalous (in the usual sense).

Insofar as plasma instability is concerned, it is important to note that oscillations were observed both in the ultrashort-wave region, corresponding to frequencies in the range 200–500 MHz, and at relatively low frequencies in the band between 10 and 45 kHz.

A high level of ultrahigh-frequency oscillations was observed only during the initial stage of the discharge, when the current through the plasma was increasing. After about 300 μsec following the onset of the discharge, the level of the oscillations was found to fall and became lower by an order of magnitude, as compared with the beginning of the hf current pulse. The skin-layer plasma density corresponding to the maximum level of ultrahigh-frequency oscillations was approximately 2×10^{13} cm^{-3}. The frequency of ion-acoustic oscillations (these experiments were carried out with helium) was found to be

$$f=\frac{\omega_{pi}}{2\pi}=\frac{1}{2\pi}\sqrt{\frac{4\pi e^2 n}{M}}=450\cdot 10^6\ \text{Hz}.$$

Since the maximum of the spectrum of ultrahigh-frequency oscillations lies precisely in this region, it may be concluded that we are dealing with ion-acoustic oscillations.

The level of these oscillations is found to fall very substantially as the discharge develops, and is relatively low in the steady state. The presence of the ion-acoustic oscillations is used by the authors of these studies to explain the somewhat too intensive heating of the ions and the dependence of the ion temperature on mass. The mean ion temperature is relatively high—of the order of some tens of electron volts—which is comparable with the electron temperature. Moreover, the heavy-gas ions are heated to a higher temperature roughly in proportion to the square root of the ion mass.

The authors consider that such rapid heating of ions and the observed dependence of the temperature on mass is due to the appearance of the ion-acoustic instability right at the beginning of the discharge development, since the observed ion heating time is much shorter than would be necessary for the heating to take place through the Coulomb collision mechanism. This heating time amounts to a few microseconds. This anomalously fast ion heating restricts the energy lifetime of the ions. As in the other high-frequency confinement systems and dynamic stabilization systems described above, here again overheating restricts the lifetime of the thermally insulated plasma. The only difference is that, in the hybrid systems, i.e., systems with a strong longitudinal quasi-constant magnetic field, there would appear to be the anomalously fast heating. This is connected with the possibility of resonance phenomena whilst, in the system described previously, the overheating time is not inconsistent with the description of the heating process in terms of the usual Coulomb collision mechanism.

The question of dynamic stabilization of MHD instabilities can also be approached from a somewhat more general standpoint. So far, we have confined our attention to the interaction of magnetic fields, for example, the field due to a direct current and some stabilizing field, either longitudinal or transverse. One of these was a high-frequency field with frequency comparable with the MHD frequency of the plasma as a whole.

In principle, it would appear that the dynamic stabilization effect will occur for a periodic variation (with the corresponding frequency) of one of the basic parameters characterizing the state of the plasma, for example, the density, pressure, change in radius, or change in the form of the cross section (see, for example, [40]).

In the final analysis, these oscillations are produced by the applied high-frequency field. However, the direct effect on instability is due to oscillations in other quantities.

As an example, consider the investigations carried out on the Toloskop installation, in which stable states were investigated in a toroidal system in the presence of a composite current (quasiconstant plus high frequency) in a quasiconstant longitudinal magnetic field.[34,56] These fields and currents were varied within the following limits: the longitudinal field H_z was varied from 0.2 to 2 kOe, the quasiconstant current from 2 to 10 kA, and the high-frequency current from 0.2 (preliminary ionization state) to 6 kA. The experiments were performed in a toroidal discharge chamber made of quartz with minor diameter equal to 6.2 cm and major diameter of 60 cm. The chamber was placed in a copper envelope with a slot and an internal diameter of 7.5 cm. The working gas was hydrogen at an initial pressure of 2×10^{-3} Torr. The mean charged-particle concentration in the plasma column, which was measured by various methods in a broad range of currents and fields, was found to be $2\times 10^{14}-3\times 10^{14}$ cm^{-3}.

When only the quasiconstant current flows through the plasma, there is a well-defined $m = 1$ instability when condition (4.3) is not satisfied. This, in general, is a trivial result. However, a transition to a stable state was observed when the constant field was combined with the high-frequency field for different values of the current modulation coefficient $\alpha = \tilde{J}/J$. The oscillations on the oscillograms showing the magnetic probe signal, the reflected signal from the 2-mm interferometer, and the photomultiplier signals recording the intensity of the H_β line and of the impurity lines are then found to disappear.

Moreover, it is found that in the stable state there is an appreciable increase in the plasma temperature. Thus, in the unstable state $T_e \approx 5-10$ eV and there are sharp changes in the diamagnetic signal which suggests rapid cooling of the plasma. In the stable state, the plasma temperature measured by the diamagnetic effect and the intensity of the spectral lines, reaches ~25–100 eV. The lifetime of the stable state amounted to several microseconds in these experiments, i.e., was roughly the same as in all the experiments described above.

Here again, this time could, of course, be associated with the overheating of the plasma and, again, one would expect that it should increase with increasing temperature, i.e., with increasing power output of the high-frequency source.

The very fact of transition to the stable state in this system was associated by the authors with radial oscillations of the plasma column due to oscillations in the magnetic pressure of the high-frequency field component of the plasma current.

However, it would be premature to draw any definite conclusion from all these data. This would, in fact, require additional experimental and theoretical data.

With this, we conclude our review of existing dynamic stabilization systems. As regards the work of our foreign colleagues, we refer the reader to a recent review by Berge [57] and the proceedings of the conference on Feedback and Dynamical Control of Plasmas,[49] where a detailed bibliography is given.

Both in the Berge review and in the conference proceedings, the main attention was paid to the theoretical analysis on the various proposals in the field of dynamic stabilization. Experimental studies in this area are clearly insufficient to enable us to draw definite conclusions. It has been shown above on the basis of existing experience that a theoretical analysis of any particular variant must be regarded as a starting point for the formulation of experiments and no more. Definite conclusions can only be based on analyses of sufficiently reliable experimental results and one must be ready to accept that, at the end of the day, very few of the original assumptions will remain.

7. CONCLUSIONS

The author hopes that the foregoing review of results and ideas will stimulate further studies of this exceedingly interesting, though not fully understood, branch of plasma physics.

One thing is clear: dynamic methods are sufficiently effective both in confinement and in stabilization systems for hot plasmas. The variants of these methods considered here, and the results of experimental studies which we have reviewed above, are sufficiently reliable to enable us to draw certain conclusions.

Dynamic methods of plasma confinement. An example of the application of this method is the traveling wave system which we have considered above. The magnetic field configuration in this system is such that the field everywhere increases in the peripheral direction, and it has been shown that the plasma can be in a state of stable equilibrium in such a field.

It is possible, in principle, that other and more effective systems, using dynamic methods of plasma confinement, will be found in the future. The use of high-frequency fields for this purpose opens up new possibilities which in practice are inaccessible for quasiconstant or constant fields. Here, one can vary the field geometry in time and thus acquire a further degree of freedom which can be controlled. Attempts must therefore be made to find solutions in which these advantages are exploited more effectively.

Insofar as systems for dynamic stabilization of plasmas are concerned, here again, we may conclude that the foregoing account clearly demonstrates the real possibility of practical devices based on these principles. In all the systems of this kind which we have considered (dynamic stabilization by longitudinal hf fields, stabilization with the aid of the hf component of the longitudinal current), there is a well-defined stabilization of MHD instabilities. In some cases, the stabilizing effect is stronger than one would expect on the basis of the criteria deduced from analyses of stationary systems.

In order to match the experimental results to the theoretical predictions, it will be necessary to carry out additional studies, both theoretical and experimental.

There have been many papers in recent years devoted to this subject and most of them were theoretical. Experimental studies are, of course, impeded by substantial technological difficulties, perhaps greater than in other branches of plasma physics. The situation is also complicated by the fact that substantial progress will only be possible when there is a substantial increase in the output power of hf systems.

It has been shown above that the plasma confinement time and the lifetime of the stable state in dynamic stabilization and high-frequency plasma confinement systems are restricted by the "overheating time." This, in its turn, is determined by the plasma temperature. The character of this dependence must be determined experimentally. If the conductivity is determined by Coulomb collisions, the "overheating time" should increase rapidly with temperature. So far, the various studies have been carried out on devices in which the temperature was too low and, therefore, the overheating time was restricted to a few microseconds. Higher temperatures were achieved with damped circuits, but here the lifetime was restricted by the damping time which was also of the order of a few microseconds.

In order to escape from this Procrustean bed, we must increase the output power of the continuously operating generator, at least by an order of magnitude, i.e., up to 500–1000 MW.

It is only then that we shall be able to establish reliably whether the lifetime is determined by overheating or the development of some kind of instability, and to determine the heating mechanism. This will also resolve the question as to whether there are anomalies in the absorption of power by plasma.

There are no fundamental difficulties in the use of existing technology for the development of such systems, but the unavoidable practical difficulties will, of course, demand considerable engineering skills before they can be overcome.

As regards further developments in this direction, with the final aim of achieving controlled thermonuclear fusion with positive energy yield, existing possibilities are, of course, inadequate. The basic problem is whether it will be possible to develop high-frequency circuits operating in the frequency band of a few megahertz with $Q = 10^7 - 10^8$. This is the basic question which distinguishes the methods considered here from other currently existing methods.

However, the rate of technological advance in this century of scientific and technolgical revolution is such that new engineering developments based on particular physical ideas frequently overtake purely physical studies. We are confident that, when the physics of the phenomena become sufficiently clear and, at the same time, it will be reliably demonstrated that one of the variants of dynamic confinement and stabilization of plasma is a realistic possibility, the necessary technologies will also become available. In this respect, the situation is hardly worse than in any of the other directions which are being researched at the present time.

The confinement of plasma by high-frequency fields may turn out to have one further very important advantage. Whilst the application of the DD reaction to the problem of thermonuclear fusion is limited by the magnetic emission by electrons, in the case of high-frequency confinement this restriction is removed because, owing to the well-defined skin effect, there is no magnetic field in most of the volume occupied by the plasma.

The DD reaction, especially in hf fields, looks at present like pure fantasy. There are enough difficulties with the DT reaction, but the situation may change in the future.

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