Magnetic plasma confinement

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A review is given of the present state of research into the physics of quasistationary confinement of high-temperature plasmas in magnetic fields. Both closed (toroidal) and open magnetic traps are considered. A major part of the review is devoted to the physics of plasma confinement in tokamaks. A brief account is given of the fundamentals of the tokamak, as well as of the problems of equilibrium, stability, cross-field heat and particle transport, and plasma heating and maintenance of current. The stellarator and the reversed-field pinch are also considered. The current status of open traps is reviewed.

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1. INTRODUCTION

It is now just over 30 years since it was first suggested that a strong magnetic field could be used for the thermal insulation of high-temperature plasma with the aim of achieving a controlled thermonuclear reaction in the plasma.[1-3] The simplicity of this physical idea, and the tempting prospect of being able to exploit the inexhaustible supply of nuclear energy stored in light elements, engaged the interest of I. V. Kurchatov who without delay set in train the organization of the necessary research. The experimental work was headed by L. A. Artsimovich. Soon after, M. A. Leontovich was brought into these investigations and, together with a small but active group of theorists, began a more fundamental development of the theory of magnetic plasma confinement. At first, the experiments were exceptionally successful: the first neutrons from a nuclear reaction in a pulsed high-current discharge in deuterium were produced[4] as early as 1952. However, careful analysis of the experimental data, performed by L. A. Artsimovich, led him to the conclusion that the neutrons were not of thermonuclear origin, but were emitted as a result of the appearance of a small group of accelerated ions.

This was the beginning of a determined effort, extending over many years, that combined a penetrating study of physical processes in high-temperature plasmas with the exploration of a great variety of different approaches to, and different configurations for, magnetic plasma confinement. The development of the necessary theory proceeded in parallel with this work.

It became increasingly clear to I. V. Kurchatov that the problem of controlled thermonuclear fusion was much more complex and difficult to master than it appeared at first. However, these difficulties merely strengthened this resolve, and he took steps to attract new research groups to the work on controlled thermonuclear fusion. In 1956, in a celebrated lecture given at Harwell,[5] I. V. Kurchatov suggested that research into controlled thermonuclear fusion should be declassified and that the efforts of the international scientific community should be combined in order to solve this problem. Step by step, a close scientific collaboration was established in research into thermonuclear fusion, and it is now hard to imagine how rapid advances can be made in this field without constant collaboration between scientists from different countries.

Many years of persistent research and fundamental investigation have ensured that the physics of high-temperature plasmas is now one of the major fields of modern physics. We now have extensive knowledge of the physical processes that occur in plasmas confined by magnetic fields. In addition to establishing definite priorities from the point of view of possible progress toward a thermonuclear reactor, contemporary research is continuing along the lines of a more detailed investigation of different principles of plasma confinement and the underlying physical phenomena. The physics of magnetic plasma confinement has, in fact, turned out to be exceedingly fruitful and interesting to both experimenters and theorists.

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2. PRINCIPLE OF MAGNETIC CONFINEMENT

The magnetic field neither accelerates nor slows down the motion of charged particles: it merely curves their trajectories. However, this effect can be very considerable. For example, a 10-keV deuteron traveling in a moderate magnetic field of, say, 30–50 kG, will move on a helix of radius amounting to a few millimeters. The trajectory of an electron in a comparable field is a helix of radius smaller by a factor of 60. It would seem, therefore, that it should be sufficient to close the magnetic tube of force into a torus in order to produce a closed magnetic trap. In actual fact, this is not the case because particles do not simply wind themselves on to the curved lines of force but, in addition, experience the so-called magnetic drift. Nevertheless, one can consider theoretically an extensive set of closed magnetic traps in which this drift will be compensated on the average.

As an example, Fig. 1 illustrates the tokamak and stellarator configurations. In the tokamak, the magnetic field is made up of the toroidal field produced by external coils, and the poloidal field (i.e., lying in meridional planes) due to the current I flowing through the plasma and the current in controlling conductors (not shown in the figure). As a result, the lines of force lie on magnetic surfaces inscribed into one another. In the stellarator, such surfaces can be produced by external conductors alone, i.e., by circular coils producing the toroidal field and special additional helical coils. In the stellarator of the torsatron type, only the helical coils are required and produce both the toroidal and poloidal field components.

Nonuniformity of the magnetic field in the longitudinal direction can be used to produce the so-called open, or mirror, magnetic traps. In the mirror traps, the particles are reflected from magnetic mirrors, i.e., regions of stronger magnetic field (this effect is related to the conservation of the transverse adiabatic invariant in the slowly-varying magnetic field). Both closed and open traps must confine not only the individual charged particles, but also the quasineutral plasma containing electric fields that appear spontaneously in the plasma. It is also desirable that plasma confinement should not deteriorate with increasing plasma density n (i.e., the number of ions per cubic centimeter) and the attendant increase in the plasma pressure $p=2nT$ (we are assuming that the electron and ion temperatures and densities are the same, and that the temperature T is measured in energy units, so that the Boltzmann constant is absent).

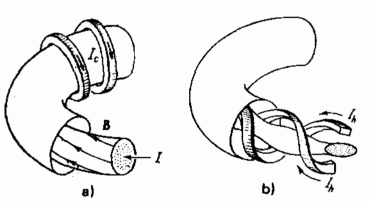

FIG. 1. Closed magnetic traps: a) tokamak, (b) stellarator (torsatron). Short arrows indicate the direction of the current: I_c in coils producing the toroidal field, I_h in helical coils, I in tokamak plasma.

In terms of the quantity $\beta=8\pi p/B^2$, which is equal to the ratio of the plasma pressure to the magnetic pressure, the above condition can be formulated as a desire to preserve the confinement properties of the magnetic trap for values of β that are not too small. The equilibrium of dense plasmas in a magnetic field can be described with adequate precision by the equations of magnetohydrodynamics. It is clear from these equations that the plasma is diamagnetic, i.e., it "pushes apart" the lines of force, so that magnetic confinement can be looked upon as the confinement of plasma by the pressure exerted by the magnetic field. Theoretically, there is a great variety of equilibrium magnetic plasma configurations. However, both theory and experiment show that many of these configurations are unstable, and theory is by no means always able to predict all the consequences of instability development. Many years of laborious experiment and careful theoretical analysis were necessary before it was possible to pick out the particular configurations that were free from the most hazardous instabilities whilst minimizing tolerable weak instabilities.

The final aim of this extensive search for satisfactory confinement of high-temperature plasmas is quite specific: most of these studies are directed toward systems in which controlled thermonuclear reactions could be achieved. This in turn means that it is necessary to produce plasmas with a temperature of about 10 keV (1 keV is equivalent to 1.16×10^7 K) and to satisfy the so-called Lawson criterion $n\tau>2\times10^{14}$ cm^{-3} s, where n is the plasma density (or more precisely, the density of ions in the deuterium-tritium mixture in which the D–T fusion reaction resulting in He^4 takes place) and τ is the characteristic time for plasma energy confinement.

Although these parameter values have not as yet been reached, the progress so far gives us every confidence to expect that they will be achieved in systems introduced in the course of the next few years. In their quest toward this tempting target, that is more applied than physical in character, the research physicists involved in this work had to enter a totally new field of physical phenomena, and to develop a description of the behavior of high-temperature fully-ionized plasmas in strong magnetic fields.

3. TOKAMAKS

(a) The tokamaks principle

The tokamak is one of the simplest and at the same time one of the most advanced ideas in magnetic confinement, both in relation to plasma parameters achieved so far and to the level of understanding of physical phenomena occurring in plasmas. Outwardly, the tokamak looks very simple: it resembles a transformer, with the secondary in the form of a closed plasma "doughnut". This plasma ring is produced by using the azimuthal electric field, induced by passing a current pulse through the primary coil-inductor, to initiate a discharge in the gas filling a toroidal vacuum chamber. The induced field also produced an increase in the plasma current to a value amounting to hundreds of thousands or even millions of amperes in large installations. A current of this magnitude will usually heat the plasma to a temperature of the order of 1 keV. Still higher temperatures are achieved by additional heating of the plasma. It is thus clear that, as far as the method of producing the plasma and maintaining the plasma ring in equilibrium is concerned, the tokamak is similar to the so-called pinched, i.e., self-compressed, discharge. The principal individual feature of the tokamak is the presence of a strong magnetic field in the toroidal direction. This field is produced by special coils and serves to stabilize the plasma.

The principle of stabilization of a current-carrying plasma column by a strong longitudinal field was proposed a relatively long time ago.[6,7] Shafranov[8,9] showed that a stable toroidal discharge can be produced in two ways, namely, either by using a very strong magnetic field, or by employing a moderate magnetic field, comparable with the field produced by the current itself. The former possibility is realized in the tokamak configuration, and the latter in the so-called field-reversed pinch (see below).

The criterion for the stabilization of plasma by a strong magnetic field is relatively simple in form:

$$q(a)=\frac{B_{\mathrm{T}}a}{B_\theta R}>1, \tag{1}$$

where B_{T} is the totoidal magnetic field, B_θ is the field due to the plasma current at the edge of the ring, and a and R are the minor and major radii of the plasma ring, respectively. The quantity q given by (1) is called the safety factor, and the condition given by (1) is referred to as the Kruskal-Shafranov condition.

This is the basic criterion for producing a stable toroidal plasma ring in the tokamak. We shall see that, for high plasma currents and the associated high values of $B_\theta=2I/ca$, this criterion reduces to the condition for a very high toroidal magnetic field. Experimentally, it is often necessary to depart quite considerably from the theoretical limit, taking $q(a)\simeq 2-3$, and it is only in the most recent and more careful experiments that it has been possible to reduce $q(a)$ to values of 1.4–1.6.

It will be clear from the foregoing why the word tokamak—an acronym for the Russian expression that translates as "toroidal chamber with magnetic coils"—has been widely adopted to describe this class of installation.

(b) Plasma equilibrium

In large tokamaks, the plasma confinement time is much greater than the characteristic time between two-body Coulomb collisions, so that, locally, the plasma is in a state approaching thermodynamic equilibrium. The particle velocity distribution is therefore nearly Maxwellian, and one can speak not only of the density $n(\rho)$ but also the local temperature $T(\rho)$, where ρ is the distance from the magnetic axis. When the electron and ion temperatures are equal, the resultant plasma pressure $p=2nT$ satisfies the equilibrium condition

$$\nabla p=\mathbf{j}\times\mathbf{B}/c \tag{2}$$

where $\mathbf{j}$ is the current density and c is the velocity of light.

It follows immediately from (2) that $\mathbf{j}\cdot\nabla p=0$ and $\mathbf{B}\cdot\nabla p=0$, so that, in equilibrium, the current and field lines lie on surfaces of constant pressure, which are therefore automatically magnetic surfaces.

The Ampere force on the right of (2) includes a contribution due to not only the product of the longitudinal current and the poloidal field, but also the product of the poloidal and toroidal fields. The presence of a strong longitudinal field thus ensures that, in contrast to the simple pinch, the pressure and the square of the poloidal field are not rigidly related. Accordingly, the quantity $\beta_\theta=8\pi\bar{p}/B_\theta^2$ ($\bar{p}$ is the average plasma pressure) can be either less than or greater than unity. For $\beta_\theta<1$, the ring has paramagnetic properties with respect to the longitudinal magnetic field, but for $\beta_\theta>1$ it behaves as a diamagnetic medium. It was doubted at one time as to whether the diamagnetic stage with $\beta_\theta>1$, in which part of the plasma pressure is maintained by the toroidal field and not the field due to the current, could be realized. However, there are now experimental data for $\beta_\theta=2$ or more, including data obtained in large installations. It is important to note that all existing experimental data on the equilibrium of plasmas in tokamaks are in complete agreement with the theory.

The equilibrium magnetic configuration in the tokamak is completely determined by the shape of the cross section of the plasma ring, which can be controlled by controlling the currents in the external "shaping" coils, and by the two functions $p(\psi)$ and $q(\psi)$ of the poloidal flux ψ, where

$$q(\psi)=\oint\frac{B_{\mathrm{T}}\,dl_{\mathrm{P}}}{2\pi rB_{\mathrm{P}}} \tag{3}$$

in which $B_{\mathrm{T}}=2F(\psi)/cr$ is the toroidal field, $B_{\mathrm{p}}=\nabla\psi/2\pi r$ is the poloidal field, and the integral is evaluated over the periphery of the cross section. The poloidal flux ψ satisfies the nonlinear differential equation[10-11]

$$r^2\,\mathrm{div}\frac{\nabla\psi}{r^2}=-16\pi^3r^2p'(\psi)-\frac{16\pi^2}{c^2}F(\psi)F'(\psi), \tag{4}$$

which is referred to as the Grad-Shafranov equation. The shape of the cross section is usually described with the aid of the following parameters: the aspect ratio $A=R/a$ (a is the transverse size of the cross section), the displacement Δ of the centre of the cross section from the axis of the chamber, the elongation K, and the "triangularity" δ of the cross section:

$$\begin{aligned} r &= R+\Delta+a\cos\theta-\delta\sin^2\theta,\\ z &= Ka\sin\theta. \end{aligned} \tag{5}$$

Efficient numerical methods have now been developed for solving (4). Thus, by investigating stability, one can determine the optimum parameters and profiles, and also calculate the necessary shaping currents.

From the point of view of maximum ratio of plasma to magnetic pressure ($\beta=8\pi p/B^2$), the D-shaped cross section with elongation $K\simeq 1.5$ and small triangularity δ is the most convenient. For this configuration, the theoretical limit for β is greater by a factor of approximately two as compared with the tokamak of circular cross section. Moreover, elongation along the vertical can be used to increase the volume available to the plasma for the same aspect ratio A, and also to increase the energy lifetime τ_E of the plasma (see below). The advantages of using an elongated cross section have been demonstrated in a number of installations, including T-8, T-9, and T-12 in the USSR, DOUBLET in the USA, and so on.

From the standpoint of controlling impurities in the plasmas, there is considerable interest in equilibrium with a separatrix for which the extreme magnetic surfaces are open and can be used to produce the so-called divertor configuration (see below, Fig. 4).

It has been shown experimentally that the plasma parameters are exceedingly sensitive to even very small shifts of the outer magnetic surface of the plasma ring relative to the chamber (or the diaphragm limiting the ring). Special steps must therefore be taken in modern tokamaks to ensure equilibrium and the precise centering of the plasma, even in the case of a circular cross section. A conducting shell was used to ensure equilibrium in the early tokamaks with pulse lengths or the order of 10 ms, but methods of controlling equilibrium without the shell were subsequently developed.[12] In the larger installations, equilibrium can be maintained for up to several seconds.

The limiting equilibrium plasma pressure is an important quantity for toroidal systems. This limit is related to the appearance of the separatrix on the plasma boundary, owing to the compensation of the poloidal field on the inner side of the torus by the external confining field $B_\perp$. The latter field can be determined from the condition for zero net force:

$$\frac{2\bar{p}\pi a^2}{R}=\frac{1}{c}IB_\perp, \tag{6}$$

and hence $B_\perp = aB_\theta\beta_\theta/2R$ where $\bar{p}\pi a^2/R$ is the tension per unit length under the influence of pressure. The factor 2 appears because of the diamagnetic expansion of the plasma ring in the toroidal field. On the inner side of the torus, the field $B_\perp$ is opposite to the field due to the current which, because of the asymmetry of the current distribution for $\beta_\theta \gg 1$, is approximately equal to $B_\theta/2$. Thus, the condition for zero resultant toroidal field on the inner side of the torus, i.e., $B_\perp \simeq B_\theta/2$, yields

$$\beta_\theta \approx \frac{R}{a}. \tag{7}$$

This formula is also valid for plasma of noncircular cross section and, since $\beta = \beta_\theta B_\theta^2/B_T^2$ and $q = aB_T/RB_\theta$, we find that the limiting value of β in a tokamak of circular cross section is given by

$$\beta \simeq \frac{a}{Rq^2}. \tag{8}$$

The limiting equilibrium pressure is sufficiently high (in a tokamak with a D-shaped cross section it is possible to reach $\beta \simeq 20$–30%) and does not give rise to any anxiety.

The restriction imposed by instability sets in before that imposed by equilibrium (see below).

(c) Neoclassical transport theory

For many years, the principal aim of tokamak research has been the determination of the laws of heat and particle transport across the magnetic field. At first sight, it would appear that diffusion and thermal diffusivity in tokamaks should be the same as in a simple direct uniform magnetic field, as long as $B_T \gg B_\theta$. However, this is not so. In current experiments, and even more so in proposed future experiments, the plasma temperature is so high that the mean free path of particles between Coulomb collisions is greater than the length of the torus. The particles therefore succeed in traversing a long portion of a line of force between collisions, so that transport is determined not by the size of the Larmor circle but by the deviation of the particles from the magnetic surfaces in the course of their drift motion. The quantative theory of transport that takes into account the actual trajectories of electrons and ions in plasma with low collision frequency was developed by Galeev and Sagdeev.[13,14] It is called the neoclassical theory.

It can be explained qualitatively as follows. Particles traveling along the lines of force can be divided into two groups, namely, those that pass right through the configuration and those that are trapped by it. Trapping occurs because the particles are reflected by the magnetic mirror on the inner part of the toroidal surface where the magnetic field is stronger. It occurs only for particles with low longitudinal velocities $v_\parallel \lesssim v\sqrt{\varepsilon} \simeq \rho/R$ and v is the resultant velocity of the particle. In the course of the oscillations between the mirrors, the trapped particles depart from the magnetic surfaces by the amount $\xi = qr_B/\sqrt{\varepsilon}$ where q is the local value of the safety factor, $r_B = v/\omega_B$ is the Larmor radius of the particles, and ω is the Larmor frequency. In the transverse plane, the trajectory of the guiding centre looks like a banana (Fig. 2).

At very low collision frequencies, it is precisely these trapped particles that provide the main contribution to the transport process, so that the thermal diffusivity can be estimated from $\chi = f\xi^2\nu_{eff}$ where ν_{eff} is the effective collision frequency and $f \simeq \sqrt{\varepsilon}$ is the fraction of trapped particles. Moreover, since distant Coulomb collisions result in diffusion-type relaxation in velocity space, we have $\nu_{eff} \simeq \nu/\varepsilon$ for trapped particles with $v_\parallel/v \simeq \sqrt{\varepsilon} \ll 1$, where ν is the ordinary collision frequency. Thus, in high-temperature collisionless plasma, the transport processes occur in the trapped "bananas" and result in thermal diffusivity of the form $\chi \simeq q^2 r_B^2 \varepsilon^{-3/2}\nu$ where r_B is the average Larmor radius. Of course, this expression is valid only if the particles do not succeed in colliding within the "banana", i.e., for $\nu_{eff} < v_\parallel/qR$ or $\nu < (v/qR)\,\varepsilon^{3/2}$. On average for all the "bananas", the velocity v must be interpreted as the thermal velocity v_T. The dimensionless quantity $\nu^* = \nu\varepsilon^{-3/2}qR/v = \varepsilon^{-3/2}qR/\lambda$ is commonly referred to as the

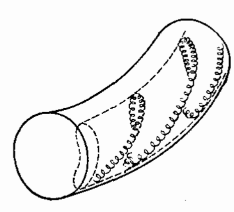

FIG. 2. Trajectories of trapped particles in a tokamak (the so-called bananas).

collision parameter. Thus, for very infrequent collisions, i.e., for $\nu^* < 1$, the thermal diffusivity can be written in the form $\chi \simeq q^2 r^2 v_T \nu^*/qR$. When $\nu^* > 1$, particles with small $v_\parallel$ succeed in undergoing a collision in a single period of their oscillation along the lines of force, so that the idea of the "banana" is no longer meaningful. The main contribution to the thermal diffusivity is then provided by passing particles that succeed in undergoing one collision in one revolution over the minor circle. Because of toroidal drift, such particles become displaced by the amount $\xi \simeq q r_B v_T / v_\parallel$, and their fraction is $f \simeq v_\parallel / v_T$. The condition that particles undergo one collision per revolution is then $\nu_{eff} \simeq v_\parallel / qR$. This leads to the following estimate for the thermal diffusivity: $\chi \simeq f \xi^2 \nu_{eff}$, and we have the so-called plateau, i.e., the diffusivity is $\chi \simeq q^2 r_B^2 v_T / qR$, and is therefore independent of the collision frequency. The plateau extends to $\nu = v_T / qR$, i.e., $\lambda = qR$ for which $\nu^* = \varepsilon^{-3/2}$. The region of frequent collisions is often referred to as the Pfirsch-Schluter region. It begins for $\nu^* \geqslant \varepsilon^{-3/2}$ for which $\chi \simeq q^2 r_B^2 \nu$. Thus, the complete expression for the thermal diffusivity is

$$\chi = q^2 r_B^2 \frac{v_T}{qR} \mathcal{F}(\nu^*), \tag{9}$$

where $\mathcal{F} \simeq \nu^*$ for $\nu^* < 1$ and $\mathcal{F} \simeq qR/\lambda$ for frequent collisions for which, $\nu^* > \varepsilon^{-3/2}$. Finally, $\mathcal{F} = 1$ in the region of the plateau. The theoretical expressions for the ion (χ_i) and electron (χ_e) thermal diffusivities are illustrated schematically in Fig. 3a. As can be seen, $\chi_i \gg \chi_e$, simply because $r_B^2 v_T$ is proportional to the square root of the particle mass, whereas the mean free paths of ions and electrons are roughly equal.

Although we have given a purely descriptive account, there are, in fact, accurate quantitative calculations of the thermal diffusivity χ. These calculations[15] provide smoother relationships for χ_i, χ_e that are free of the corners shown in the figure. As far as the plasma diffusion coefficient D is concerned, this coefficient is of the order of χ_e in neoclassical theory, i.e., it should be substantially smaller than χ_i. The neoclassical theory also enables us to calculate the diffusion of impurities relative to the main plasma, and suggests the possibility of other very interesting transport phenomena.

One of them is the so-called trapped-particle pinch which was originally considered qualitatively by Ware.[16] This arises for the following reason. When trapped particles are reflected by the magnetic mirrors as they approach the inner surface of the torus along the lines of force, they describe "banana" type trajectories that do not run around the magnetic axis. When the induced electric field E appears along the torus, they begin to drift with a velocity v_d in the radial direction, and this velocity is determined by the equilibrium between the Lorentz force $e v_d B_\theta / c$ and eE. Since the fraction of trapped particles is approximately equal to $\sqrt{\varepsilon}$, the resultant macroscopic velocity turns out to be $v_d \simeq (cE/B_\theta)\sqrt{\varepsilon}$. The neoclassical theory enables us to find a more accurate value for v_d throughout the range of collision frequencies, including the weaker passing-particle pinch effect.

Another effect that is symmetric with respect to pinching is often referred to as the bootstrap current (this current is maintained by the plasma itself and the effect is similar to that mentioned in the Alice in Wonderland story in which she holds herself up by her own shoelaces—hence the name "bootstrap"). This effect occurs for $\beta_\theta > \sqrt{R/a}$ when the poloidal β for the trapped particles, i.e., $\beta_\theta \sqrt{a/R}$, is greater than unity.[17,18] When this is so, the trapped particles can push apart the poloidal lines of force in the course of their diffusion, thus preventing the decay of the field due to finite conductivity. The theory enables us to calculate this effect precisely (it has not been observed in tokamaks and stellarators, possibly because of the presence of small magnetic-field fluctuations[19]). The neoclassical theory also enables us to calculate many other fine effects such as the effect of particle trapping on the longitudinal electrical conductivity, the viscosity, the effect of field corrugation on transport, and so on. In brief, the neoclassical theory is now complete, and the question is whether experimental data do or do not agree with this theory.

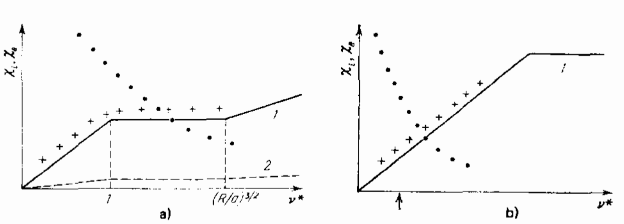

FIG. 3. Thermal diffusivity of plasma: (a) current experiments, (b) future experiments on tokamak reactors. Arrow shows the working point. Solid line—theoretical χ_i, crosses—experimental χ_i (qualitative), dashed and dotted curves—theoretical and experimental (qualitative) χ_e. $\nu^* = (qR/\lambda) \times (R/a)^{3/2}$; λ is the mean free path.

(d) Thermal conductivity of plasmas

Studies of the thermal conductivity of plasmas across the magnetic field have been, and continue to be, at the center of attention in experiments on tokamaks, since it is precisely this conductivity that is the main indicator of the efficiency of magnetic plasma confinement. The first attempt to interpret experimental results in this way in a systematic fashion was undertaken by Artsimovich.[20,21] He showed that, in the few experiments that had been carried out at the time, the loss of heat along the ion channel was in satisfactory agreement with the neoclassical theory in the region of the plateau (the experiments were performed in this region), but heat transfer along the electron channel was much greater than the theoretical prediction. New tokamaks were subsequently built, so that the range of plasma parameter values achieved in different installations was considerably extended, and more complete data were obtained on the ion and electron conductivities.

Above all, the use of additional (as compared with the Joule) heating resulted in considerably higher plasma temperatures, so that the region of infrequent collisions could be reached (for which the transport of trapped-particles on the "bananas" becomes important). Plasmas with $\nu^* = \nu q R / v_T \varepsilon^{3/2} \ll 1$ were produced in the T-11 and PLT installations, where v_T is the thermal velocity and $\varepsilon = a/R$ is the toroidality. The region of infrequent collisions right up to $\nu^* \simeq 10^{-2}$ was particularly well penetrated in the PLT for which ion temperatures $T_i(0) \simeq 7$ keV were achieved for $n_e(0) = 4.5 \times 10^{13}$ cm^{-3}. The results of these experiments are in good agreement with the neoclassical theory, i.e., to within a factor of 2 or 3. This is illustrated schematically in Fig. 3a where the crosses indicate that experiment is in good agreement with theory both in the region of the plateau and in the region of infrequent collisions (these crosses are not experimental points but merely a schematic representation of the experimental result).

Thus, there is good agreement between theory and experiment for the ion channel, and we may conclude that heat transfer along the ion channel in future thermonuclear reactors will be small provided only it is not specially enhanced, for example, in order to control the rate of the D–T reaction (ion heat transfer can be enhanced by a slight corrugation of the magnetic field, which may appreciably affect the trapped particles).

Extensive experimental data are now available on the electron channel as well. By varying the plasma parameters in different installations, particularly through additional heating, and by comparing the results obtained for different installations, attempts were made to find purely empirical relationships or, as they have come to be known, scaling functions, for the energy confinement time τ_{E_e} in the electron channel, or the electron thermal diffusivity χ_e. Different workers proposed different empirical formulas for τ_{E_e} and χ_e, which were often quite different from one another but represented experimental results with comparable precision. This was connected, on the one hand, with the fact that the experimental data themselves were subject to considerable spread and, on the other hand, with the fact that ohmic heating introduced an additional relationship between the plasma parameters (the ohmic heating power can be expressed in terms of the same quantities as χ_e), which complicated the determination of the pure relationships. It is only very recently that further experiments, using additional heating, have made the situation much clearer.

Nevertheless, we would like to mention the so-called Alcator scaling which, at one time, was used to systematize a particular range of experimental data on χ_e in tokamaks with ohmic heating. This was, indeed, one of the simplest and, therefore, quite popular scaling function obtained on the Alcator-A installation in which the density could be varied within very broad limits by varying the magnetic field. As a result, it was shown that it is precisely the plasma density and not its other parameters, such as temperature, magnetic field, and current, that influences the plasma energy confinement time. In particular, the energy confinement time in this installation turned out to be simply proportional to the plasma density. Comparison with experimental data obtained for other installations led to an empirical formula that was subsequently still further simplified by the design group working on the INTOR tokamak reactor.[22] This simplified formula is

$$\tau_E = 5\cdot 10^{-19}\bar{n}a^2\,(\text{c}). \tag{10}$$

It is sometimes referred to as INTOR scaling ($\bar{n}$ is the average plasma density in cm^{-3} and a is the minor radius of the torus in cm). The simplified Alcator scaling, i.e., the INTOR scaling, corresponds to very simple empirical expressions for the thermal diffusivity ($\chi_e = 5\cdot 10^{17}/n$ cm^2 s^{-1}) and thermal conductivity ($\varkappa_e = n\chi_e = \text{const} = 5\times 10^{17}$ cm^{-1} s^{-1}). To a very crude approximation, these expressions are in reasonable agreement with a wide range of experimental data, especially those on Joule heating. However, they suffer from a major defect in that they are purely empirical. They are written in terms of plasma parameters, so that there is no guarantee that they will remain valid when they are extrapolated to the future larger installations. By analogy for the procedures adopted in other branches of physics, for example, in hydrodynamics, in which empirical laws in the form of dimensionless groupings of parameters are employed, it seems natural to use dimensional analysis to obtain an expression for χ_e.[23,24]

It is convenient to start with the expression for χ_e in terms of the other natural physical parameters of the tokamak plasma. Quite recently Mukhovatov, Merezhkin, *et al.*[25,26] carried out a series of very accurate experiments on the T-11 installation with diaphragms to investigate χ_e as a function of ρ, and to obtain an empirical expression that can be written in the form

$$\chi_e \cong \frac{\varepsilon^{7/4}}{qRnr_e}\left(\frac{2T}{mA}\right)^{1/2} \tag{11}$$

where $\varepsilon = \rho/R$, ρ is the distance of a given point from the magnetic axis, q is the local safety factor, R is the major radius of the torus, $r_e = e^2/mc^2$ is the classical radius of the electron, T is the local electron temperature, m is the mass of the electron, and A is the atomic mass of the ion.

The Mukhovatov-Merezhkin empirical formula is in reasonable agreement with other experimental data and, in particular, with the measured local values of χ_e in the T-10 and FT installations.[27,28] It may well be that the predicted dependence on ε is too strong: recent experiments with the Alcator-C installation[29] are in better agreement with a linear dependence on ε.

As can be seen, the expression given by (11) shows that χ_e does in fact vary with the density as n^{-1}, whereas the dependence on the other parameters is much weaker because, as a rule, the temperature in the larger installations is higher and the parameters a/R and q vary within a relatively narrow range.

Equation (11) has a definite theoretical justification.[30,31] In particular, it is clear that it can be rewritten so that it contains the dimensional grouping $(v_e/qR)c^2/\omega_{pe}^2$, where v_e is the electron thermal velocity and $\omega_{pe} = \sqrt{4\pi e^2 n/m}$ is the Langmuir frequency. The ratio c/ω_{pe} corresponds to the well-known depth of penetration of the magnetic field into the collisionless electron gas (in superconductors this is the London field penetration depth), and qR/v_e is the characteristic time of electron transit along the magnetic field. Accordingly, the electron thermal conductivity can be related to the so called "magnetic flutter"[32-35], i.e., small oscillations of the magnetic surfaces due to plasma noise. As they approach one another, the magnetic surfaces can "exchange" lines of force as a result of reconnection over a width $\delta \sim c/\omega_{pe}$. As a result, the lines of force are weakly stochasticized with line-of-force diffusion coefficient $D_B \sim \delta^2/qR$, where qR is the characteristic length over which reconnection takes place. The field stochasticity leads to the transport of electrons with thermal diffusivity $\chi_e = v_e D_B$. The additional dependence on ε appears because the main contribution to reconnection is due to the field-perturbation component on the outer circuit.

We thus see that equation (11) is essentially collisionless in character and contains no dependence on the collision parameter ν^*. Recent experiments,[36-40] involving strong additional heating in T-11, ASDEX, DOUBLET-III, ISX-B, and PDX, have resulted in data on the thermal conductivity of plasmas with high values of β_θ and β—the ratio of plasma pressure to the pressure due to the poloidal and toroidal magnetic fields, respectively. The record value $\beta = 4.7\%$ was reached in the DOUBLET-III system for sufficiently acceptable plasma energy confinement, and somewhat lower values of β have been obtained in other systems. Some deterioration in plasma confinement is observed in all the experiments with high values of β_θ, β, and this cannot be explained exclusively in terms of a deterioration in the spatial distribution of the energy due to additional heating as compared with ohmic heating. The deterioration in confinement can be represented by an additional factor F in (11). So far, this factor does not exceed 2–3 even in the ISX-B and PDX installations for which the highest deviations have been observed from the "soft limit" type ohmic-heating scaling in β_θ (in ISX-B) or in $q\beta$ (in PDX). The factor F reflects the increase in χ_e with increasing β_θ or $q\beta$, so that it includes additional dependence on dimensionless parameters that are not present in (11). It also turns out that F has a favorable dependence on ν^*, i.e., it decreases with ν^*. Thus, the fact that plasma energy confinement deteriorates with increasing β_θ, β need not, as yet, be viewed pessimistically from the standpoint of extrapolation to the larger future installations. It is more likely that current experiments reflect the presence of ballooning modes[41,42] whose importance will decline with increasing plasma temperature.

Thus, the final picture that emerges is as illustrated schematically in Fig. 3. To within a factor of 2–3, the ion thermal conductivity agrees with neoclassical theory and will be low in future experiments (Fig. 3b). The electron thermal conductivity is anomalous, i.e., it exceeds the neoclassical value by 1 or 2 orders of magnitude, and has a completely different dependence on the plasma parameters (for example, it is independent of ν^* for small β_θ). It is qualitatively clear from Fig. 3b that, in future installations, it will be the electron channel that will be largely responsible for plasma energy loses. The empirical result indicated by (11) for χ_e is fully acceptable for future thermonuclear installations, but a more precise elucidation of the possible enhancement of this quantity for high β is an urgent problem.

(e) Plasma diffusion

Thermal conductivity is not the only transport process that takes place in a tokamak plasma. There is also particle transport across the magnetic field, i.e., diffusion of the main hydrogen plasma and of the impurities, redistribution of the current and of the poloidal magnetic field, viscous redistribution of the angular momentum, and transport of neutral gas due to charge transfer in the plasma. All these processes are important for the understanding of the physical phenomena occuring in plasmas, and there is now an extensive volume of experimental data relating to these processes, although the resulting picture is neither clear nor complete. This refers in particular to plasma diffusion.

From the standpoint of classical theory, there is a considerable difference between diffusion of pure hydrogen plasma and diffusion of impurities in hydrogen plasma. The point is that, since pure hydrogen plasma is quasineutral, the ion current must be exactly equal to the electron current, and this is achieved by readjusting the radial electric field that appears when the ambipolarity of the diffusion process is slightly disturbed. Hence, according to the neoclassical theory, the coefficient of diffusion of hydrogen plasma turns out to be of the order of the electron thermal diffusivity χ_e, i.e., very small.

However, the electron thermal conductivity is actually highly anomalous, so that we cannot expect the actual diffusion coefficient to be close to the theoretical result. Numerous experiments have, in fact, shown that the measured plasma diffusion coefficient is anomalous and is usually[43] between $0.2\chi_e$ and $0.3\chi_e$. The diffusion process is accompanied by pinching at a rate approaching the neoclassical value at very high densities. However, experiments performed on the ALCATOR-A, FT, and PLT systems have shown that, when the plasma density is very high, there is enhanced pinching, accompanied by enhanced diffusion on the periphery. This problem is being actively investigated.

Insofar as the usually multiply-charged impurities are concerned, their diffusion is not closely related to the electron component because the electrical charge transported by impurities can be compensated by a reverse current of hydrogen. Since the ion thermal conductivity can also be looked upon as a countercurrent of ions of different energy, the impurity diffusion coefficient D_Z must, clearly, be of the order of χ_i, i.e., even in the theory of binary collisions, it must be substantially greater than the theoretical χ_e. Next, since an impurity ion carrying Z charges replaces Z hydrogen ions, the clustering of multiply-charged ions at the centre of the discharge is convenient from the standpoint of minimum chemical potential, since it can give rise to the radial expansion of the hydrogen component. This effect is, in fact, predicted by classical theory, although there is also the reverse effect of impurity-ion repulsion due to thermal diffusion which is proportional to the temperature gradient.

The experimental picture of the behavior of impurities is quite complex because the transport phenomena predicted by the neoclassical theory occur against the background of anomalous diffusion of both the main plasma and of the impurities, and both processes take place. One can therefore occasionally observe effects associated with increased impurity density at the centre but, as a rule, when noise is at its usual moderate level, the diffusion coefficients of impurities with different Z are close to one another, and both impurity clustering at the center and thermal-diffusion effects are important. There are indications that the impurity diffusion coefficients decrease with increasing atomic number of the main plasma.[29,44] A number of methods for directly influencing the impurity flux have now been proposed and are designed to extract the impurities from the plasma. This can be done by using particle beams, resonance hf waves, or asymmetric gas supply on the periphery. All these methods have been the subject of extensive experimental investigation.[40,42,45-47]

(f) Impurities

Impurities play an important role in tokamaks. Even low impurity levels will lead to considerable radiation loss, largely in the form of line emission which is itself unfavorable from the point of view of plasma confinement. However, in addition, impurities have a further uncontrollable effect on the radial distribution of temperature and current density, leading to the formation of profiles that are subject to MHD instability. It was precisely the development of methods for suppressing impurity levels in plasmas that has led to advances in achieving the desired parameters in tokamaks. Moreover, impurities may also play a positive role in the large future installations, since the radial temperature distribution can be varied in a desired way by suitably choosing the impurity concentration and composition.

Typical impurities consists of "light" components such as carbon and oxygen, which are present in great abundance on the surface of the vacuum chamber, and "heavy" components, consisting of atoms of the material of the walls and limiting apertures, and produced by evaporation under the impact of hot ions and neutrals. Light impurities are useful in modern installations, since they are completely ionized in the central region and do not contribute to the line emission whereas, on the periphery, they lead to the cooling of the plasma and to a reduction in the rate of arrival of heavy impurities. In reactor plasma, the light-impurity concentration can be of the order of a few per cent, whereas the concentration of heavy impurities should not exceed a few hundredths of a per cent.

Effective methods have now been devised for freeing plasmas from impurities, so that sufficiently pure plasma with $Z_{eff} \simeq 1$ can be produced. In future large-scale installations, in which the plasma temperature and the wall load will be higher, the impurity problem will become more acute and will require special attention. The most effective way of protecting the plasma from heavy impurities is to use the so-called divertor, i.e., a special magnetic-field configuration capable of diverting all or some of the current from the surface

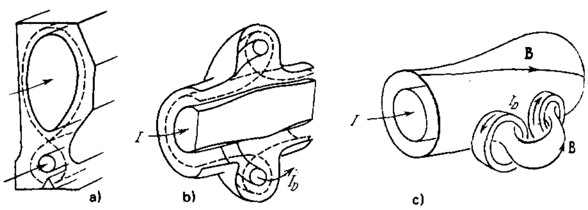

FIG. 4. Divertors: (a) poloidal with D-shaped plasma cross section, (b) toroidal, (c) bundle.

of the plasma along the lines of force into special volumes in which the interaction of the plasma with the wall (or more precisely with the target plates) can be reduced. Examples of divertor configurations are shown in Fig. 4.

The efficiency of divertors has been examined experimentally on installations of intermediate size, such as DIVA, T-12, and DITE,[48-50] and on large installations such as DOUBLET-III, ASDEX, and PDX[51,52,37]. These experiments showed that the use of the divertor gives rise to a substantial (roughly by an order of magnitude) reduction in the concentration of heavy impurities in the plasma. A particularly interesting divertor with a high cold-plasma density in the divertor chamber was investigated on the DOUBLET-III and ASDEX installations. It was shown that a kind of gas cushion (plasma blanket) is produced in this case and received the plasma current from the main chamber, redistributing it over the large surface of the divertor chamber. The use of the divertor not only reduces the interaction between the plasma and the walls, but also provides a way of removing the products of the thermonuclear reaction (α particles), which is essential for the steady-state operation of a reactor.

(g) MHD instability of plasmas

The tokamak plasma is a nonequilibrium thermodynamic system. It acts as a large reservoir of free energy: magnetic energy associated with the current and gas, and kinetic energy associated with the pressure. Both these energy components can become sources of large-scale magnetohydrodynamic instabilities, producing a rearrangement, or even the destruction, of the plasma configuration. In tokamaks with $\beta_\theta \lesssim 1$, the main reservoir of free energy is associated with the current necessary to maintain the toroidal equilibrium.

The concept of resonance perturbation is of major importance in the analysis of instabilities in toroidal systems. This is a helical perturbation of the form $\exp i(m\theta - n\varphi)$ whose pitch is equal to the average pitch of a magnetic line of force. The condition for this to happen is

$$nq = m. \tag{12}$$

A magnetic surface on which this condition is satisfied is called a resonance surface. Resonance perturbations do not produce local bending of magnetic lines of force, which requires the expenditure of energy. By acting along the entire magnetic line of force, even a small resonance perturbation is capable of transporting the line of force to a large distance away from its unperturbed position. If the pitch of the lines of force on different surfaces is not constant, i.e., field shear involving the crossing of the lines of force is present, so that $s = -\rho q'(\rho)/q \neq 0$, then the perturbation will change sign along the line of force at a certain distance from the resonance magnetic surface, and its overall effect will be reduced. Shear will thus limit the width of the perturbed layer. The structure of the layer depends on the electrical conductivity of the medium. In vacuum, or in a medium with finite conductivity, the structure of the magnetic surfaces in the perturbed layer will assume an island configuration. The resonance surface will thus effectively split into m helical filaments. This splitting is topologically forbidden in the approximation of perfect conductivity, and the resonance surface exhibits singularities: functions and their derivatives are discontinuous, and a surface current appears.

The largest scale instability is the so-called screw instability in plasmas with a free boundary for which the conductivity is zero outside the current channel. Screw instabilities can occur when the resonance magnetic surface lies outside the current channel or near its boundary. They produce a deformation of the surface of the plasma ring, and affect its central part as well. The curvature of the torus plays no essential role in this process and, in the theoretical analysis of the situation, the torus can be replaced by a cylinder with identically defined ends. For small deformations of the cylinder, the displacement vector of a volume element of the plasma can be represented by a set of harmonics of the form $\xi = \xi(\rho)\exp i(m\theta + n\varphi + \gamma t)$. Theoretical analysis of the instability of the current channel began with the most obvious kink instability ($m = 1$ mode). It is precisely for the stabilization of this mode that the use of the strong longitudinal magnetic field was suggested, and this field eventually became one of the main elements of the tokamak. The stability condition (1) for the $m = 1$ mode, i.e., $q(a) = aB_T/RB_\theta(a) > 1$, signifies that a magnetic line of force on the boundary of the current channel will rotate during its passage along the entire torus through an azimuthal angle less than 2π (since otherwise the screw pitch will be greater than the length of the system).

The longitudal magnetic field will stabilize the $m = 1$ mode but it can give rise to an instability for modes with $m \geqslant 2$, which deform the cross section of the plasma ring without affecting the shape of its axis. Early studies of instability with idealized current density profiles led to the pessimistic conclusion that the screw instability was unavoidable for one of the higher modes. Fortunately, more detailed analysis[53] showed that the instability set in only for certain ranges of values of $q(a)$, namely,

$$m - \alpha < nq(a) < m, \tag{13}$$

where the coefficient α and the growth rate γ for modes with $m \geqslant 2$ are very dependent on the current profile (Fig. 5). It turns out that $\alpha \leqslant 1$ for the uniform and surface currents, and the instability zones over the entire range of values of $q(a)$. However, for the bell-shaped current-density profile, the instability zones contract ($\alpha < 1$) and stability windows appear between the instability zones. These windows are what makes a tokamak viable.

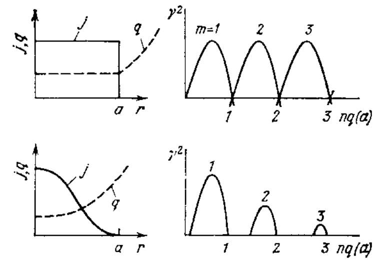

FIG. 5. Stability windows in the tokamak for helical perturbations of the form $\exp[i(m\theta - nz/L)]$: (a) slab current profile, (b) bell-shaped profile.

The presence of the successive instability zones can be explained as follows. To the left of the left-hand boundary of an instability zone, the deformation of the plasma ring is prevented by the magnetic field B_θ due to the current itself. To the right of this boundary, the curved current-carrying ring experiences the destabilizing effect of the longitudinal field. As the perturbation length decreases, there is an increasing contribution due to the excitation of the surface current (which is actually distributed within a thin layer) that accompanies the displacement of the highly-conducting plasma relative to the longitudinal field. When $nq(a) > m$, its stabilizing effect exceeds the destabilization due to the interaction between the curved current and the external field.[54]

The screw instability is most clearly evident in the case of a current confined to the skin-layer, which corresponds to a nonmonotonic profile $q(\rho)$ with a minimum in the interior of the plasma. This instability prevents the skin-layer current distribution from developing. As the current grows in the tokamak, one observes bursts of oscillations as $q(a)$ passes through integer values (between $q > 10$ and $q \sim 2$–3), and the bell-shaped current profile is established. This phenomenon has been simulated numerically,[55] using the connection between the initial and final states, established for processes involving the reconnection of magnetic surfaces.[56] The instability condition given by (13) corresponds to the situation where the resonance magnetic surface falls into the interval $a < \rho < \sqrt{m/(m-\alpha)}\,a$ outside the current channel. The development of the screw modes is accompanied by the splitting of this resonance magnetic surface.

When the resonance surface $q(\rho) = m/n$ lies in the interior of the ring, the screw mode cannot develop in the approximation of perfectly conducting plasma because the condition for a frozen-in magnetic field is restricted by the demand that the field must not change its topology. However, a slower dissipative screw mode, referred to as the tearing mode, can develop under these conditions. The development of this mode is associated with the reconnection of magnetic surfaces and the formation of their filamentary structure. The criterion for the appearance of the tearing instability if qualitatively the same as for the ideal screw instability. The instability will develop when there is a sharp drop in the current density and the resonance surface lies outside the surface of maximum gradient. The mode becomes stabilized when the resonance surface enters the interior of the main current channel. This is well known to experimenters. Rapid cooling of the edge of stable plasma (on contact with a diaphragm, or during injection of neutral gas, accompanied by radiative cooling of the periphery) leads to the contraction of the current channel. Stability is found to deteriorate rapidly if, at the same time, the resonance surface lies outside the current channel.

It has been shown theoretically that the bell-shaped current profile with a small gradient in the surface layer can ensure stability against large scale (low m, n) screw and tearing modes. It is becoming increasingly clear that experimental techniques for producing stable modes, which require special preparation of chamber walls in order to avoid uncontrolled influx of impurities, careful control of equilibrium, programed current growth, and injection of neutral gas into the discharge, are equivalent to the establishment of the optimum current profile for stability. The utilization of all this technology has resulted in the realization of stable modes for $q(a) < 2$ instead of the $q(a) \simeq 3-5$ in experiments performed in the 1960s. This means that the current was increased without increasing the longitudinal magnetic field, with the attendant improvement in plasma heating and confinement.

Even when the conditions for macroscopic stability are satisfied in the plasma, internal MHD instabilities can still develop. One of them,[53] namely, the internal screw mode with $m=1$, is also associated with the current and not the pressure. It arises when the value of q on the axis becomes less than unity, as a rapidly-varying current profile, which is favorable for the stabilization of the above surface modes, becomes established. The central part, which is bounded by the resonance surface $q(\rho_s)=1$, can then curve along a screw. In the transverse section, the inner part is shifted outward and is pressed against the resonance surface. This is a weak mode in the case of perfect conductivity, but the introduction of dissipative effects, leading to the reconnection of magnetic surfaces at the point of their contraction, results in a considerable increase in the growth rate. This instability is readily detected experimentally in the form of saw-tooth temperature oscillations in the central region (see below).

In addition to the large-scale screw instabilities that are present in the tokamak and appear even at zero plasma pressure, there is a further possible instability due to the plasma pressure gradient which is common to all magnetic traps. This instability is connected with the displacement of the plasma-containing magnetic tubes of force to a position of lower total energy, and can develop in the neighborhood of closed magnetic lines of force. The region in which it is localized is restricted by shear, i.e., a change in the pitch of the lines of force. In the absence of shear, the instability corresponds to the Rayleigh-Taylor instability of a liquid located in the gravitational field if its density increases with height. In the case of magnetic confinement, the curvature of the magnetic lines of force is the analog of the gravitational force. When the lines of force are convex and shear is absent, the plasma is unstable.

Moderate-pressure plasma is stabilized by shear $[q'(\rho)\neq 0]$ and also by a variation in the sign of the normal curvature of the magnetic lines of force which, taking (2) into account, is given by

$$k_n = -\mathbf{n}\cdot\frac{\mathbf{B}}{B}\nabla\frac{\mathbf{B}}{B} = -\frac{\mathbf{n}\cdot\nabla(B^2+8\pi p)}{2B^2}, \tag{14}$$

where n is the outward normal to the magnetic surface. On the outer side of the current we always have $k_n>0$, whereas on the inner side the line of force is concave for $q>1$. Because of the high (sonic) velocity with which pressure equalization occurs along a magnetic tube of force, the average curvature of a line of force that is "seen" by plasma with $\beta_\theta \lesssim 1$ turns out to be negative. Since positive curvature corresponds to a reduction and negative curvature to an increase in the magnetic field away from the above surface, the possibility of a negative average curvature is equivalent to the possibility of an average magnetic well. Because of the apperance of the magnetic well in the tokamak, the criterion for local stability of moderate-pressure plasma, $\beta_\theta \lesssim 1$, is very favorable:[57]

$$\frac{s^2}{4} + \frac{8\pi p'(\rho)\rho}{B^2}(1-q^2) > 0. \tag{15}$$

Stability is assured when, locally, $q(\rho)>1$. The magnetic well will also stabilize dissipative local modes that are not influenced by shear.

All the above types of instability are associated with moderate-pressure plasma, i.e., $\beta_\theta \lesssim 1$, which corresponds to $\beta=\beta_\theta(a/Rq)^2\sim(a/Rq)^2$. Calculations show that, for reactors, we must have $\beta \gtrsim 5\%$. As already noted, such values of β present no problem from the point of view of equilibrium, but can influence stability. The point is that, when $\beta_\theta \sim R/a$, the gas-kinetic pressure is sufficient to produce a curvature of magnetic tube of force on the outer circuit around the torus when the ends of the tube are fixed on the inner circuit in the region of negative curvature. The pressure for which this ballooning instability sets in can be calculated by the elegant method developed by Connor *et al.*,[58] for studying small-scale modes ($n\to\infty$), which requires the use of two-dimensional equilibrium calculations.[59] The low ballooning modes are investigated directly by solving numerically the eigenvalue problem for the two-dimensional linearized equations of motion.[60] For plasmas of circular cross section, the limiting value of $\bar{\beta}$ turns out to be about 2.5%.

In the tokamak in which the plasma ring has a *D*-shaped cross section, which is the most favourable shape from the point of view of exploiting the negative-curvature effect (magnetic well), this limit rises to about 5%. Calculations lead to the following interpolation formula (which can be qualitatively justified) for $\bar{\beta}_{cr}$:

$$\bar{\beta}_{cr} \approx \frac{b}{5qR}; \tag{16}$$

where b is the vertical semiaxis of the transverse cross section of the plasma. These results were obtained for $q(\rho)$ between $q(0) \simeq 1$ and $q(a) \simeq 2$ for typical values $R/a \simeq 4$–5. Existing experimental data do not exclude the possibility of modes with $q(0) < 1$. In that case, the limiting pressure for stability turns out to be higher and is entirely sufficient from the standpoint of reactor requirements.

Because of the anomalous electron thermal conductivity, experiments with ohmic heating alone cannot produce β_θ greater than unity. The use of additional heating (mainly injection of neutral atoms) has resulted in values $\beta_\theta \sim 0.5R/a$, which is close to the theoretical limit, and $\bar{\beta} \simeq 3\%$ in tokamaks with circular cross section of the plasma ring, for example, in T-11. The record value $\bar{\beta} = 4.7\%$ has recently been obtained in the DOUBLET-III tokamak (with a D-shaped plasma) for semiaxial ratio $K = b/a = 1.4$. The fact that such values can be reached with currently available heating-power levels is due to the exploitation of regimes with low $q(a)$ (< 2). It has been shown experimentally that an increase in the heating power leads to the saturation of β. The corresponding limit is referred to as "soft" because it does not lead to catastrophic consequences, and merely results in the enhancement of transport which, as noted in subsection (e), is related to the stochastization of magnetic lines of force under the influence of the ballooning instability.

(h) Disruptive instabilities

MHD oscillations in the form of small-amplitude helical modes are almost always present in tokamaks. They are usually observed with the aid of small coils (the so-called Mirnov coils) placed outside the plasma ring. As the amplitude of these perturbations increases, one also observes "magnetic islands" within the transverse sections through the magnetic surfaces.

In addition to this low-level MHD activity, the tokamak plasma is occasionally also subject to a very menacing phenomenon that is referred to as disruptive instability. This instability is accompanied by a loss of energy from the plasma and attendant strong MHD oscillations, as well as the radial compression of the current density distribution. The disruptive instability can lead to the complete breakup of the plasma ring and to the termination of the current in plasma, in which case one speaks of a "major disruption". There are also minor disruptions, in which energy is lost from the plasma, and which recur without resulting in the complete disintegration of the plasma ring. In minor and, even more so, in major disruptions, the circuit-voltage oscillogram $V(t)$ shows the presence of negative bursts that correspond to the expulsion from the plasma of a fraction of the poloidal magnetic field.

The disruptive instability was for a long time a complete puzzle, and the key to its understanding emerged only after another interesting phenomenon, namely, the so-called sawtooth oscillations, was discovered.[61] The basic phenomenon involved here is x-ray emission by the central part of the plasma ring. It was shown (initially on the PLT and subsequently on many other tokamaks) that soft x-rays that are emitted from the

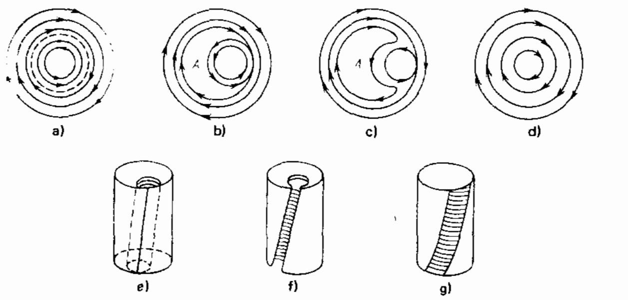

FIG. 6. Reconnection of magnetic surfaces in the axial region of a tokamak: (a–d) transverse cross section, (e–g) segment of torus.

centre of the ring and are a measure of the central temperature T_e, exhibit sawtooth oscillations in which each "tooth" is several miliseconds long and constitutes a slowly-growing intensity with a sharp fall to a level of the order of 10% at the end of the tooth. This fall is accompanied by a sharp reduction in the temperature at the centre. Studies of the spatial distribution of these oscillations have shown that the reduction in temperature occurs only within a certain region $\rho < \rho_s$ in the interior of the ring, i.e., ρ_s is usually appreciably less than a. Outside the region defined by ρ_s, the "teeth" change sign, and the slow fall in temperature is followed by a sharp rise. More detailed experiments also show that the $m = 1$ oscillations are excited for $\rho < \rho_s$ prior to the temperature drop at the center.

All these "internal disruption" phenomena in the sawtooth oscillations can be naturally explained by the development of an internal $m = 1$ mode for $q < 1$ ($q = 1$ occurs precisely for $\rho = \rho_s$), followed by a reconnection of the lines of force.[56,62] How this occurs is illustrated in Fig. 6. It is assumed that, initially, the quantity q is less than unity at the centre of the ring and rises toward the periphery, becoming equal to unity on the broken curve in Fig. 6a. Relative to the $m = 1, n = 1$, helical perturbation, the lines of force have a positive transverse component inside the region $\rho = \rho_s$, and a negative component outside this region, i.e., the pitch of the lines of force is less than $2\pi R$ in the interior of the $\rho = \rho_s$ surface, and greater than this value elsewhere. When the helical disturbance appears, the lines of force cross (Figs. 6b and c) and then reconnect (Figs. 6c and d). The inner part of the ring is thus pushed out (Figs. 6e–g). A more precise picture of the reconnection process in sawtooth oscillations emerged as a result of numerical simulation.[62,63] Experimental data are consistent with this picture although there are indications of a more complex "turbulization" process during the ejection of heat out of the $\rho < \rho_s$ region.

A similar, though more complicated, picture of reconnection of lines of force is also found to emerge in the course of the disruptive instability which usually begins with the development of a helical disturbance with $m = 2, n = 1$. When this disturbance and the corresponding magnetic island are large enough, interaction with other modes (e.g., the internal mode with $m = 1, n = 1$) can begin and gives rise to new modes and the corresponding islands. When these islands come into contact with one another, the lines of force may reconnect, and the current distribution may become compressed in the radial direction.[64] In the final analysis, the energy source for this instability is the energy of the poloidal magnetic field. This is why the process can be understood in terms of the idealized picture involving the penetration into the interior of the plasma of helical "vacuum bubbles",[65] and can be interpreted in terms of the rupture of the ring due to the interaction between the current and the longitudinal field.[54] Disruptive instability can be avoided by controlling the radial profile of the current distribution.

(i) Kinetic instabilities

In addition to large-scale instabilities that are treated within the framework of magnetohydrodynamics, less hazardous small-scale instabilities may develop in the tokamak plasma. They reflect the effects of the finite Larmor radius and of the differences between the motion of different groups of particles. They can be analyzed theoretically within the framework of the transport equation with self-consistent electric and magnetic fields, and are therefore referred to as kinetic instabilities. They involve, above all, different types of drift instability associated with the excitation of drift waves. It seemed at one time that these instabilities were universal and unavoidable even in the idealized case of a plane plasma layer and, even more so, in the more realistic toroidal geometry. Actually, drift oscillations in a plane layer are found to be stable when shear is taken into account.[66-69] Factors associated with the growth of small-scale oscillations appear in the toroidal geometry, especially when trapped particles are taken into account.[70] These instabilities present a definite theoretical hazard for magnetic plasma confinement. In reality, however, the density and temperature fluctuations observed in tokamaks have not as yet been directly related to the anomalous electron thermal conductivity and diffusion. It seems that drift-type instabilities lead not so much to transport as to perturbations of the magnetic surfaces, producing slight disruptions. The resulting anomalous transport depends less strongly on the energy source for the magnetic fluctuations, i.e., on the shape and level of drift-type oscillations, at least for plasmas with β_θ, β well away from critical values.

Kinetic-type instabilities include oscillations excited by runaway electrons which are occasionally observed in tokamaks. We recall that runaway electrons are the electrons in the tail of the Maxwellian distribution which, having acquired energy in the longitudinal electric field, achieve a "breakthrough", i.e., cannot be retarded because the Coulomb cross section falls rapidly with energy. The accelerated electrons can excite waves through the so-called anomalous Doppler effect: an electron gyrating in the magnetic field is similar to a rapidly moving oscillator, so that a wave whose phase velocity in the frame in which the electron is at rest is less than the velocity of the electron appears as a wave of negative energy. It follows that the electron can excite a wave by increasing its transverse and reducing its longitudinal velocity components.[71] Bursts of oscillations accompanied by the ejection of fast electrons appear in the tokamak plasma as a result of the excitation of this type of "fan" instability. The theory of this phenomenon is in very good agreement with experiment.[72-75]

Kinetic instabilities will undoubtedly become the subject of fundamental theoretical and experimental research, since they reveal a new range of very interesting collective plasma phenomena. From the practical point of view, their effect on the thermal insulation of the plasma is not very severe.

(j) Additional heating and maintenance of the current

The necessity for additional plasma heating as a way of reaching thermonuclear temperatures was recognized at an early stage of research into controlled fusion. This includes, above all, high frequency methods of plasma heating in different wavelength ranges, and the corresponding experiments were conducted over many years. However, it became clear in the mid-seventies that the injection of fast neutral atoms is preeminent among these methods. It was first used on the CLEO tokamak at Culham and was subsequently widely adopted because of the simplicity of the physical principle underlying ionization trapping of fast particles in plasmas, their good confinement, and also the rapid advances in the techniques available for producing well-collimated high-intensity atomic beams. The injection of fast atoms was in fact used to reach the highest ion temperatures of about 7 keV in the PLT and the highest value of $\bar{\beta}$, i.e., 4.7%, in DOUBLET-III. The injection method is also promising from the point of view of initiating the thermonuclear reaction in future thermonuclear reactors. However, the technique is not very simple and requires further development before it can be extended to the higher beam energies (up to 150–200 keV) that are necessary in the larger installations.

High-frequency and microwave methods of plasma heating have therefore attracted increasing attention in recent years and new important data have been obtained as a result. The greatest advances have been achieved in experiments on the heating of plasmas by fast magnetoacoustic waves at frequencies close to the ion cyclotron frequency and its harmonics. The greatest increase in the temperature (to about 3 keV), was achieved in the PLT installation, using a small amount of ^3He impurity in a deuterium plasma.[76] It was shown that the heating of this small admixture was similar to the injection of fast atoms, so that the associated physical phenomena were also similar.

Studies have also been carried out of the heating of plasmas at higher frequencies, namely, in the region of the electron cyclotron and hybrid (intermediate between electron and ion cyclotron frequencies) resonances. Heating at the electron cyclotron frequency is particularly effective.[77] Thus, in recent experiments performed on the T-10 installation, using 0.5 Mw of microwave heating power, the electron temperature was raised to about 3 keV for a density of 4×10^{13}–5×10^{13} cm^{-3}. This method produces less disturbance in the plasma. Moreover, by suitably choosing the position of the resonance, it is possible to influence the temperature and current-density distributions and, hence, the stability of the plasma. The disadvantage of the method is the complexity of the microwave sources of gyrotron waves. Much effort is therefore being devoted to another range of lower-hybrid waves that is more acceptable from the engineering point of view. The physics of the excitation and absorption of lower-hybrid waves is more complex and the experimental data are not as yet sufficient to enable us to estimate the future possibilities of this method of heating.

If the plasma is heated by some external means, the problem naturally arises as to whether the current in the plasma can also be maintained by external means. Under ordinary conditions, friction between the electron and ion gases is compensated by means of an induced vortical electric field. However, the current can be maintained by this inductive method for only the limited period of time while the current is growing in the inductor. But momentum can be transferred to the electrons not by the induced vortical field but by beams or by high-frequency waves. The first experiment on maintaining the current in this way was performed on the DITE installation. Using tangential injection of a beam of fast atoms into the plasma ring a current of about 30 kA was produced with the total current being equal to 100 kA. The possibility of generating a current by means of Alfven waves has been demonstrated on a small toroidal system at the Sukhumi Physicotechnical Institute.[78] Recently, striking results were obtained with high-frequency waves close to the lower-hybrid resonance. As an example, Fig. 7 shows circuit voltage and current oscillograms for a power level of about 200 kW and sufficiently low plasma density of less than 10^{13} cm^{-3} on the T-7 tokamak with superconducting magnetic field coils.[79] As can be seen, the circuit voltage falls to zero during the high-frequency pulse, i.e., the current in wholly maintained by the high-frequency field. Currents is excess of 200 kA have also been produced in PLT and ALCATOR-C.[76,80] The density in the latter installation was quite high, namely, $n = 4\times10^{13}$–6×10^{13} cm^{-3}. The theoretical efficiency with which the current can be maintained at the lower-hybrid resonance falls with increasing density, but is still about 0.1 A/W in thermonuclear plasma at a temperature of 10 keV and density of 10^{14} cm^{-3}, i.e., this figure is fully acceptable for a stationary tokamak reactor. The experimental efficiency is, so far, much lower than the theoretical figure, and additional experimental work will be needed on the development of this method of maintaining the current.

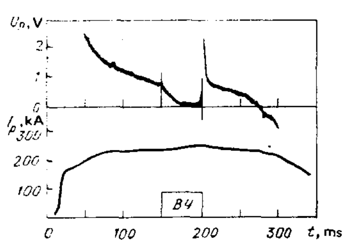

FIG. 7. Maintenance of current by a high-frequency field in the T-7 tokamak.

4. STELLARATORS

(a) Physical fundamentals

The stellarator is a device in which a system of inscribed toroidal (in the topological sense) magnetic surfaces is produced by external currents.[81] In other words, the line integral of **B** over a poloidal contour drawn around any magnetic surface in the stellarator must be equal to zero:

$$\oint \mathbf{B}\cdot d\mathbf{l}_p = 0. \tag{17}$$

At first sight this condition is inconsistent with the fact that a magnetic line of force is wound on a toroidal magnetic surface in such a way that the vector **B** must have a regular poloidal component. The resolution of this difficulty lies in the somewhat unusual helical geometry of the stellarator magnetic surfaces. In the so-called "ordinary" stellarator with a circular axis, the transverse cross section of the magnetic surfaces is not circular but in the form of an ellipse (or a rounded triangle, square, etc.) that rotates as one moves along the system. The vacuum magnetic field in this type of stellarator has both toroidal and helical components. If we neglect the curvature of the torus, we can derive this field from the potential ($\mathbf{B} = \nabla\phi$)

$$\Phi = B_0\Big[R\varphi + \sum_l \varepsilon_l \frac{Rl}{m} I_l\left(\frac{m\rho}{R}\right)\sin(l\theta - m\varphi)\Big]; \tag{18}$$

where B_0 is the field on the axis of radius R, ρ is the polar radius, θ, φ are the angle variables, and ε_l is the relative amplitude of the helical harmonic with l and m periods along the minor and major circuits around the torus, respectively. For the sake of simplicity, the number m is assumed to be the same for all the harmonics. If the dominant harmonic in (18) is that with $l = 2$, we speak of a two-bulge stellarator or, simply, an $l = 2$ stellarator. Since the equations $\mathrm{div}\,\mathbf{B} = 0$, $\mathrm{curl}\,\mathbf{B} = 0$ are identical to the equations describing the irrotational flow of incompressible liquids, the behavior of the magnetic lines of force on a magnetic surface can be visualized by considering flow-lines in an ideal incompressible liquid flowing in a tube having the same shape as the magnetic surface. In a tube in which the cross section rotates as we move along the axis, the current lines are also set (purely kinematically) in rotation and, on average, turn through an angle $\overline{\Delta\theta} = \mu_1$ in one period of the helix. This angle increases as the cross section departs from circular shape. If, for example, the cross section has a corner, the entrainement on the corresponding helical line of the tube is complete, and $\mu_1 = 2\pi$. The local entrainement velocity has a maximum at the apices of the elipse (or the "triangle" etc. as the case may be). On compressed segments, the current line not only lags behind the rotation, but also moves for a certain time in the reverse direction. It follows that, if at the apices of a noncircular cross section the magnetic line of force is inclined in the direction of rotation, then on compressed segments of the magnetic surface it is inclined in the opposite direction. The resulting alternation of the sign of the poloidal component of **B** explains the validity of the condition given by (17) and the systematic rotation of the line of force in the azimuthal direction. The equation of a magnetic line of force on a magnetic surface can, in general, be written in the form

$$\theta = \theta_0 + \mu\varphi + \theta_1(\theta_0, \varphi), \tag{19}$$

where $\theta_1(\theta_0, \varphi)$ is a periodic function that depends on the parameter θ_0 characterizing the chosen line of force. The quantity μ is the average ratio of the number of revolutions of a line of force in azimuth θ to the number of revolutions along the torus

$$\mu = N\mu_1 = \lim_{\varphi\to\infty} (\theta/\varphi), \tag{20}$$

and is referred to as the rotational transform ($N = m/l$ is the number of complete revolutions on the helical surface). On a surface of elliptic cross section with semiaxes a, b, we have $\mu_1 = (a-b)^2/(a^2+b^2)$. As the cross section tends to the circular shape, $a \to b$ and the rotational transform decreases. For example, it is zero on the axis of the three-turn ($l = 3$) stellarator.

When the rotational transform is not zero, the separation of charges that occurs as a result of the inhomogeniety of the toroidal magnetic field (the so-called toroidal drift) is removed by the current $\mathbf{j}_\parallel$ flowing along the magnetic lines of force. Its distribution is determined from the continuity equation $\mathrm{div}(\mathbf{j}_\parallel + \mathbf{j}_\perp) = 0$ where $\mathbf{j}_\perp$ is the diamagnetic current that is related to ∇p by the equilibrium equation (2)

$$\mathbf{j}_\perp = c\frac{[\mathbf{B}\cdot\nabla p]}{B^2}, \quad \mathbf{B}\cdot\nabla\left(\frac{j_\parallel}{B}\right) = -[\mathbf{B}\cdot\nabla p]\cdot\nabla\frac{1}{B^2}. \tag{21}$$

The main component of the secondary (relative to the diamagnetic) current $\mathbf{j}_\parallel$, which is equal to the Pfirsch-Schluter current, is of dipole origin and has opposite directions on the inner and outer parts of the section of the torus:

$$j_\varphi \cong \frac{2}{cB_T}\frac{dp}{d\rho}\cos\theta. \tag{22}$$

The above relationships are also valid for tokamaks in which the rotational transform $\mu = 1/q$ is produced by the current I excited in the plasma. The main difference, as compared with the tokamak, is that, in the stellarator, the effective poloidal magnetic field $B_\theta^{st} \simeq (\rho/R)B_T\mu$ is external. This is the origin of the importance of the Pfirsch-Schluter currents in the balance of forces. In contrast to the tokamak in which the plasma ring is prevented from expanding in accordance with (6) by the external transverse field $B_\perp$ that interacts with the current I, in the stellaratory it is held in position by the external poloidal field B_θ^{st} during its interaction with the φ component (22) of the Pfirsch-Schluter currents.

(b) Equilibrium and stability of plasmas

The above difference between the stellarator and the tokamak is very favorable from the point of view of plasma stability. Thus, firstly, the absence of the ohmic heating current ensures that the stellarator is free from large-scale helical instabilities or disruptive instabilities, which are particularly troublesome in the tokamak. Secondly, the external system of helical fields and the corresponding external poloidal field give the magnetic configuration a definite rigidity. Formally, this is reflected in the fact that the rotational transform and the magnetic well, which govern plasma stability, depend not only on the derivative of the radial displacement $\Delta'(\rho)$, but also on the absolute value of $\Delta(\rho)$. The equation for the latter quantity in the stellarator is[82,83]

$$[(\mu_{st}\Delta)'\rho^3]' + \frac{1}{\mu_I}(\mu_I^2\rho^3\Delta')' = -\frac{8\pi p'(\rho)\rho^2 R}{(\mu_{st}+\mu_I)B_0^2}; \tag{23}$$

where μ_{st} is the rotational transform which depends on the geometry only, and is due to the current. The solution of the homogeneous equation ($p' = 0$) corresponds to the imposition of a uniform transverse magnetic field. As can be seen, in the tokamak ($\mu_{st} = 0$), the uniform transverse field shifts the plasma ring as a whole without changing the relative position of the magnetic surfaces, and $\Delta'(\rho) = 0$. On the other hand, in the stellarator with shear $[\mu_I = 0, \mu_{st}'(\rho) \neq 0]$, the solution of the homogeneous equation $\Delta(\rho) = \mathrm{const}/\mu_{st}(\rho)$ shows that the disposition of the magnetic surfaces is altered. In the typical case of positive shear, the inner magnetic surfaces are shifted more than the outer surfaces and, if the shift is directed away from the center of curvature of the axis, an average magnetic well is formed (the average magnetic field on the outer magnetic surface is stronger than on the axis which is shifted so that the toroidal magnetic field is reduced). Thus, the imposition of a uniform transverse magnetic field can be used in the stellarator to vary the depth of the magnetic well within certain limits. This depth is important for the stable confinement in the stellarator of plasmas at sufficiently high pressure.

Since the magnetic well is due to the inhomogeneity of the toroidal magnetic field, it can be produced only for a relatively low aspect ratio. The corresponding limiting (with respect to stability) value of the parameter β_{stab} decreases with increasing aspect ratio R/a (Fig. 8). According to (23), the limiting equilibrium β, determined from the condition $\Delta \simeq a$ (which roughly corresponds to the appearance of the separatrix on the plasma boundary), is given in accordance with (23) by the same expression as in the tokamak:

$$\beta_{eq} \approx \frac{a}{R}\mu^2. \tag{24'}$$

In contrast to the tokamak, the rotational transform is not limited by large-scale instabilities. For a given pitch of the helical surfaces, it increases with R/a. Correspondingly, β_{eq} increases with the aspect ratio. The intersection of the $\beta_{stab}(A)$ and $\beta_{eq}(A)$ curves determines the optimum values of the aspect ratio and the limiting plasma pressure. Calculations show that, typically, $A_{opt} \simeq 10$ and $\beta \simeq 5$–10%.

The limiting pressure turns out to be substantially

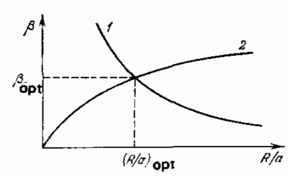

FIG. 8. Qualitative dependence of the equilibrium and stability limited parameter $\beta = 8\pi\bar{p}/B^2$ on the aspect ratio R/a in an $l = 2$ stellarator: 1—stability limit, 2—equilibrium limit.

higher in the stellarator with a helical axis, but such stellarators are constructionally more complex.

Detailed studies of plasma equilibrium and stability in stellarators are based on numerical computations. In contrast to the tokamak—a system with axial symmetry—we no longer have an exact scalar equation of equilibrium of the form of (4). The magnetic field is now an essentially local characteristic, so that the stellarator—a three-dimensional system of inscribed magnetic surfaces—is described in term of the special language of flux variables. Instead of the components of the vectors $\mathbf{j}=(c/4\pi)\,\mathrm{curl}\,\mathbf{B}$ and $\mathbf{B}$, it is usual to introduce integrated quantities such as the toroidal current I and flux ϕ within a given magnetic surface, and poloidal current F and magnetic flux ψ outside this surface. The required quantities are the functions $r(a,\theta,\zeta)$ and $z(a,\theta,\zeta)$ describing the transformation to the curvilinear flux coordinates a,θ,ζ, where a is a marker on the magnetic surface and θ,ζ are arbitrarily chosen cyclic coordinates in the toroidal geometry. The resulting system of scalar equations is still too complicated but, if we suppose that the helical amplitude is small, it can be reduced to a single equation of the form given by (4), and the computational methods developed for tokamaks can be employed again. This procedure has been successfully used to perform quite detailed studies of equilibrium and stability of plasmas in stellarators.

(c) Studies of plasma confinement

Despite the obvious attraction of stellarators (time-independent magnetic system, absence of current driven instabilities), they have not been as widely adopted as tokamaks. One of the reasons for this can be seen in the disappointing plasma confinement results at the Princeton laboratory where stellarator studies have been in progress since the original proposal by L. Spitzer in 1951. Experiments performed in the 1960s on Stellarator C at that laboratory showed the persistent presence of the so-called Bohm confinement-time relationship $\tau\propto B/T$, which is unfavorable from the point of view of thermonuclear confinement (weak dependence on magnetic field, strong deterioration in confinement with increasing temperature). In the 1970s, under the influence of the successful experiments with the T-3 tokamak at the Kurchatov Institute of Atomic Energy, in which plasma temperatures of about 1 keV were attained, Stellarator C was converted into Tokamak ST, and thereafter the Princeton Laboratory did not return to experiments with stellarators. It is now considered that the reason for the lack of success with Stellarator C—the first relatively large stellarator with helical coils—was that the magnetic field was not precise enough. The point is that in stellarators, which are essentially three-dimensional systems, the harmonics of the helical magnetic field can combine with curvature effects to give rise to resonance perturbations of the form $\exp i(m\theta-n\varphi)$, which "split" the so-called rational magnetic surfaces with $\mu=n/m$ into m helical filaments. The island structure of magnetic surfaces then appears in the transverse sections. The size of these islands increases with decreasing shear $s=\rho\mu'/\mu$. When the amplitude of the resonance perturbations is such that neighboring islands must intersect, it is found that stochastization occurs in the perturbed tubular layer, i.e., the magnetic lines of force exhibit random drift.[85] When this extends to the entire volume occupied by the plasma, theoretical estimates show that Bohm diffusion should occur.

In the Soviet Union, Kurchatov's original plan provided for extensive research into stellarators at the Khar'kov Physicotechnical Institute. At the same time, M. S. Rabinovich initiated and directed research in this field at the Lebedev Physics Institute of the USSR Academy of Sciences. From the very outset, these studies were characterized by a very careful approach to the production of magnetic surfaces. As a result, in 1965, it was demonstrated on the L-1 installation ($l=2$ stellarator with a circular axis) at the Lebedev Institute that the confinement of plasma with $T\approx 10-15$ keV was better by an order of magnitude than in the case of Bohm diffusion[86]. At the same time, classical diffusion was observed in cesium plasma ($T\simeq 0.2$ keV) in Stellarator W-1 at the Max Planck Institute.

In 1972, ion-cyclotron heating was used on the Uragan installation at the Khar'kov Physicotechnical Institute to produce plasma with ion temperature of a few hundreds of eV, and its energy confinement time τ_E was determined experimentally.[87] It turned out to be close to the neoclassical value (in the region of the "plateau"). However, systematic studies of the current-free regime could not be carried out at the time. The first-generation stellarators had a small plasma radius ($a\lesssim 5$ cm) and a considerable length. This meant that it was difficult to produce high-temperature plasmas with low impurity levels. Second-generation stellarators were built in the 1970s at a number of laboratories across the world, and these are better suited to the use of ohmic heating. These are the relatively small ($R\simeq 1$ m, $a\simeq 10$ cm) Stellarators L-2 at the Lebedev Institute, CLEO at Culham, and JIPP T-2 at Nagoya University, as well as the larger installations WVII-A at the Max Planck Institute, HLEIOTRON-E at the Kyoto University, and Uragan-3 at the Khar'kov Physicotechnical Institute (the last machine has just been put into operation. Experiments with ohmic heating have resulted in plasmas with parameters similar to those attained in tokamaks. Anomalous energy losses via the electron channel have been shown to be reduced when the current is reduced. However, extrapolation to the case of current-free plasma turned out to be impossible. As the current was reduced to zero, the power available was insufficient to maintain the plasma.

Systematic studies of current-free plasma confinement began at the beginning of the present decade. Injection of 27-keV neutral hydrogen atoms into deuterium plasma was used in the WVII-A installation with a simultaneous reduction of the ohmic heating current to produce current-free plasma with sufficiently high parameters: $\bar{\beta}=0.5\%, n=7\times10^{13}$ cm^{-3}, $T_i=1$ keV, $T_e=0.7$ keV, $\tau_E\approx 10$ ms. The plasma was maintained for the duration of injection[88]. In the large stellarator HELIOTRON-E ($R=220$ cm, $a=20$ cm), plasma was produced by electron-cyclotron heating alone, using a pulse length of 40 ms. The plasma parameters were: $T_e(0) = 1.1$ keV, $T_i(0) = 120$ keV for $\bar{n}_e = 5 \times 10^{12}$ cm^{-3}. The plasma decay time after the end of the microwave heating pulse was 10 ms[89].

It is important to note that these somewhat optimistic data were obtained in the region of the "plateau" of the neoclassical transport theory. In the region of infrequent collisions, the neoclassical theory predicts an enhancement of transport in stellarators (in contrast to tokamaks for which it predicts a reduction), due to the trapping of particles moving along the magnetic lines of force by inhomogeneities of the helical magnetic field. Particles trapped in the helical field, drift with velocity $v_{\nabla B}$ in the toroidal field, and are displaced by the amount $\Delta_h = v_{\nabla B}/\nu_{\text{eff}}$. Hence it follows that the diffusion coefficient is $D \simeq f\Delta_h^2 \nu_{\text{eff}}$, where the fraction f of trapped particles increases with decreasing ν_{eff} so long as the collision frequency ν_{eff} remains greater than the frequency of precession around the magnetic axis. The overall transport picture is complicated by the great variety of charged-particle trajectories in the stellarator[90]. Analytic calculations performed on the assumption that the transverse dimensions of the trapped-particle orbits were small, have led to a somewhat pessimistic conclusion about the thermal insulation of thermonuclear plasmas. However, numerical Monte Carlo computations, capable of taking into account the finite nature of the orbits[91-94], have shown that the actual transport picture may well be more favorable. It is also possible to optimize the magnetic configuration with the aim of reducing the deviation of the drift orbits from the magnetic surfaces, and even controling transport by using high-frequency power to release trapped particles. Studies of transport processes in the region of infrequent collisions hold the key to future experiments on stellarators.

Many research centers are currently intensifying theoretical work on the optimization of stellarator-type systems. This includes the development of modular coil systems for the stellarator. The idea of using discrete coils of noncircular cross section to produce the stellarator magnetic configuration was first put forward and implemented on the Tor-2 system at the Lebedev Institute[95,96]. Several laboratories are developing modular systems incorporating slightly twisted coils[97]. At the same time, work is proceeding on systems with optimized Pfirsch-Schluter currents (W-VII AS)[98], systems incorporating straight segments that have only a slight effect on equilibrium, Drakon-type systems[99], systems with a novel design of the helical torus (HELIAX)[100], and so on.

5. PARAMAGNETIC PINCHES

At the beginning of 1958, English physicists made the intriguing announcement[101] that they succeeded in using the large toroidal installation ZETA to produce plasmas with a density of about 5×10^{13} cm^{-3} and the then unprecedented temperature of about 400 eV. The ZETA installation relied on the stabilization of the discharge by a weak longitudinal magnetic field trapped inside a conducting plasma. The observed paramagnetic profile of the longitudinal field, with negligible field outside the current channel, appeared to be in accord with theoretical ideas on the trapping of the "frozen-in" longitudinal magnetic flux by plasma compressed as a result of the pinch effect. On the initiative of I. V. Kurchatov, it was decided to verify these results. With the energy and speed so characteristic of Igor Vasil'evich, the work at NIIÉFA was reorganized and, in less than six months, the Al'fa installation with parameters similar to those of ZETA was built. It soon became clear that the plasma produced in ZETA was in fact cold ($T \simeq 50$ eV). The magnetic-field profile was uniquely determined by the external conditions, as was particularly clearly demonstrated by the experiments of Babichev *et al.*[102] who used a small machine built at the beginning of that year at the Kurchatov Institute. The only possible explanation was that based on quasistationary plasma turbulence.[103] The observed level of magnetic-field fluctuations was very high—of the order of 10% or more. The disappointment was so great that the unique (even in comparison with modern installations) ZETA system, and soon after the Al'fa machine as well, were dismantled.

However, an important discovery was made before ZETA was finally dismantled: it was found that there were quiescent discharge regimes in which the fluctuation level was much lower and the temperature rose to 150 eV. It turned out that, during this phase, the longitudinal magnetic field outside the current channel was negative, i.e., reversed. Since the resultant magnetic flux inside the conducting chamber could be regarded as conserved, the appearance of the reversed field had to be interpreted as evidence for the generation of additional magnetic flux in the interior of the plasma. Theoretical analysis showed[104] that, when the field is reversed, the plasma can be stable against the development of the fastest ideal MHD modes up to $\beta \simeq 30\%$, and against tearing modes up to $\beta \simeq 17\%$. The necessary condition for stability is that $q(r) = \rho B_T(\rho)/RB_\theta(\rho)$ is monotonic which, for a decreasing $B_T(\rho)$ profile, is possible only if $q(\rho)$ is inverted. The subsequent history of paramagnetic pinches, i.e., pinches with a reversed magnetic field, was somewhat tortuous. The initial attempt was to introduce active reversal of the magnetic field with a high time rate of change of the direction of the field. These experiments were performed on relatively small systems ($a \simeq 5$–7 cm), with the current flowing for periods of the order of tens of microseconds, and showed that the discharge relaxed to a stable state in a relatively turbulent way. The plasma acquired considerable amounts of impurity and remained cold. The next series of larger installations ($a \simeq 10$–25 cm) was then used to demonstrate the presence of quiescent regimes with self-reversal that were distinguished by a lower level of fluctuations ($\delta B/B \simeq 1\%$). The guiding idea in the development of our understanding of processes involving these pinches with a reversed field is due to Taylor[105] who pointed out that the reversed-field state corresponds to minimum energy for given toroidal flux and given integral

$$K = \frac{1}{c} \int \mathbf{A} \cdot \mathbf{B} \, dV, \tag{25}$$

where **A** is the vector potential ($\mathbf{B} = \mathrm{curl}\,\mathbf{A}$). This integral takes globally into account the condition of high plasma conductivity ($\mathrm{d}K/\mathrm{d}t = -2\int(\mathbf{E}\cdot\mathbf{B})\,\mathrm{d}V = -2\int(\mathbf{J}\cdot\mathbf{B}/\sigma)\mathrm{d}V \to 0$ as $\sigma \to \infty$). The condition $K = \mathrm{const}$ does not, however, restrict small-scale reconnections of the magnetic lines of force during the relaxation process.

Minimization of energy under the above restrictions leads to the relation $\mathbf{j} = \alpha\mathbf{B}$ with $\alpha = \mathrm{const}$, which corresponds in a cylinder to the solution $B_z = B_0 J_0(\alpha\rho)$, $B_\theta = B_0 J_1(\alpha\rho)$ with the possibility of negative B_z on the periphery of the column. For a relatively slow variation of the current in the discharge, its parameters succeed in relaxing to the state of minimum energy and fit well into the Taylor diagram which gives $\theta = B_\theta(a)/\overline{B}_z$ as a function of $F = B_z(a)/\overline{B}_z$.

It was feared for some time that the magnetic profile produced in the system would relax to the uniform distribution with the skin-effect time constant, so that hot plasma would not be confined for the long periods of time necessary for a reactor. To obviate this difficulty, Okawa[106] suggested that the reversed field should be maintained by external helical fields whose effect on the current should result in a "translational" transformation of the poloidal flux into a torroidal flux, mainly outside the current channel However, it was established on a number of installations (in Italy, Japan, England, and the USA) that another striking effect was taking place, namely, that of the magnetic dynamo, i.e., the continuous generation of a toroidal magnetic flux in the plasma, with the necessary energy drawn from the poloidal magnetic field. The most striking results were obtained on the ZT-40(M) installation at Los Alamos[107]. By using external circuits to maintain the constancy of the toroidal flux in the interior of the chamber, and by controlling the position of equilibrium of the toroidal plasma ring, it was possible to maintain the quiescent discharge for 20 ms or more, which is much longer than the time necessary for relaxation to the uniform distribution. Since the mechanism responsible for the generation of the magnetic flux is of turbulent origin, it enhances transport processes and results in values of τ_E that are lower than in tokamaks. It is, however, possible to show that, as the plasma temperature and electrical conductivity increase, the power necessary to maintain the paramagnetic magnetic-field distribution will diminish and the quasistationary paramagnetic pinch can also be used as a basis for the thermonuclear reactor.

Paramagnetic pinches have recently led to a new area of research concerned with the "spheromak". This system is based on a compact ($R = a$) paramagnetic configuration without the toroidal-field coils. The plasma torus can then be transported along the axis of a cylindrical chamber and, for example, a cyclically operated reactor can be achieved. Several ways of producing the spheromak configuration have been proposed and approved, but experimental work is still at a very early stage.

Greater advances have been made with compact toruses without any toroidal field, whose origins can be traced back to the θ pinches and the open high-β traps. As the axial field grows in the plasma-containing cylinder in the reversed field, the internal and external lines of force close into ellipses elongated in the direction of the axis. The resulting compact toroidal configuration is surprisingly stable (as was demonstrated by Kurtmullaev *et al.*), but this is not as yet fully understood. The plasma parameters reached with these compact systems are relatively high. Thus, in Tor-2 and FRX-C, the ion temperature reaches[108,109] 1–1.5 keV for $n = 3\times10^{15}$–4×10^{15} cm^{-3}. The observed relatively slow instabilities connected with the rotation of the plasma can be stabilized by applying quadrupole magnetic fields.

Finally, the straight z-pinches that were used in 1952 to observe for the first time the fusion reaction in deuterium plasma, have led to the investigation of the "plasma focus", "micropinches", and other systems related to pulsed or inertial controlled fusion.

6. MIRROR TRAPS

Mirror, or open, traps incorporating magnetic plugs constitute the second (after the tokamaks) area in which advances have been made in relation to controlled thermonuclear fusion with magnetic confinement. They are attractive because of their relatively simple magnetic field geometry (with straight field axis), and are subject to relatively undemanding restrictions imposed by the equilibrium conditions (essentially confined to $\beta < 1$). Moreover, the thermonuclear combustion products can be readily removed from such systems through the ends of the trap.

The principle of axial confinement of plasmas in mirror traps is based on the conservation of the transverse adiabatic invariant $mv_\perp^2/B$ during the motion of charged particles in a strong magnetic field. It has been shown experimentally that this invariant is conserved to a high degree of precision. The only unavoidable losses are those of particles with a small ratio of transverse to axial velocity (loss cone). However, initially, these investigations had to be confined to plasma instabilities and means of suppressing them. Successful solution of this problem has now brought to the fore the problem of improving particle confinement in such traps. In fact, the entire history of mirror traps is essentially the history of how difficulties in the physics of magnetic confinement have been overcome. It has revealed potential improvements that could be introduced into thermonuclear systems and high-temperature plasma confinement.

The experiments of M. S. Ioffe *et al.*,[111] have shown that the MHD-stability of plasmas in mirror traps can be assured by using minimum-B fields. It has also been shown that the so-called drift-cone instability, which is particularly hazardous for dense plasmas, can also be stabilized.[112,113] Hot plasmas can be stabilized by suppressing the inversion of the particle-energy distribution that is associated with the loss cone and occurs in the region of the characteristic values of the phase velocity of perturbations of the drift-cone instability. To achieve this, it is sufficient to add a few per-

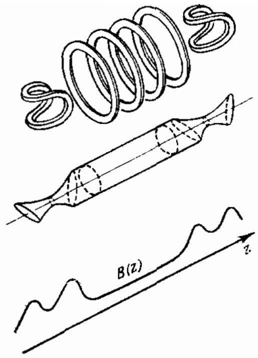

FIG. 9. The ambipolar trap: the field coils, the magnetic surface, and the field profile.

cent of "warm" plasma with a temperature of 5–10% of the temperature of the hot plasma. The best parameter values have been reached on the 2XIIB machine at the Lawrence Livermore Laboratory[114,115] ($n = 10^{14}$ cm^{-3}, $T_i = 12$ keV, $T_e = 100$ eV, $\beta \simeq 1$, $\tau_B \simeq 10^{-3}$ s).

However, even in completely stable plasmas, the leakage of particles through the mirrors as a result of Coulomb scattering is too high, and acceptable power gain ($Q \simeq 5$) cannot be achieved in the simplest versions of the trap. A possible way out of this difficulty was suggested by G. I. Dimov[115] in the USSR and by T. K. Fowler and B. G. Logan[117] in the USA. They proposed the ambipolar or tandem trap in which the confinement of the main plasma is improved by adding a further trap at each end of the main trap. The additional traps have a stronger magnetic field and contain plasma at a higher temperature as compared with the central trap. The higher positive potential produced in these additional traps, $\varphi \simeq (T_e/e)\ln(n/n_0)$, confines the ions to the central trap. The minimum-B field that stabilizes the whole plasma need only be established in the end traps (cf. Fig. 9).

The small system Gamma-6 at Tsukuba in Japan was initially used to demonstrate the possibility of ion confinement by the above potential distribution in the plasma of the ambipolar trap. Experiments performed on the TMX trap at Livermore and, later, on its upgraded version, the TMX-U, have confirmed improved confinement of plasma with β up to 30% in the central trap.[118]

Acceptable plasma parameters turned out to be very difficult to attain in the original version of the ambipolar trap, and therefore much hope is now being invested in another idea for reducing axial energy losses from plasma. This is the so-called thermal barrier.[119] Thus, whilst the principle of ambipolarity is directed toward reducing ion losses, the thermal barrier is intended to reduce the demands imposed on the plasma parameters by controlling axial electron transport. In particular, it is proposed that additional separating axisymmetric mirror traps with a moderate magnetic field be inserted between the central trap and the end traps. These intermediate traps should be filled with "bouncing ions", i.e., ions that are reflected from the magnetic mirrors near the magnetic-field maxima. Trapped ions accumulating near the field minimum in the separating traps can be removed, for example, by converting them into "bouncing ions" through charge transfer onto injected atoms. This can be used to produce a density minimum at the center of the separating trap, and the electron clouds in the central and separating traps can be completely separated by introducing a microwave-heated cloud of hot electrons into this field minimum. The presence of the thermal barrier for electrons can be exploited to increase the temperature of electrons in the end traps, and to set up the same confining potential under less stringent conditions for the plasma in these traps.

The fact that density distribution with a valley can be produced, and that "bouncing ions" can be employed in the above way, has already been demonstrated experimentally on the TMX-U. Experiments with thermal barriers in which the electron-cyclotron resonance will be used to heat electrons are in the course of preparation. This topic will undoubtedly have to be investigated experimentally, since there are theoretical doubts about the feasibility of the thermal barrier because of possible beam instabilities associated with passing ions that are highly anisotropic in the lower-density region.

The ideas of ambipolarity and of the thermal barrier are directed toward reducing longitudinal losses. However, a reduction in longitudinal losses naturally leads to the consideration of transverse diffusion processes as well.

Moreover, the utilization of quadrupolar end traps, ensuring the minimum-B configuration, may result in the enhancement of transverse transport because particle orbits then become more tortuous[120,121]. This phenomenon is similar to neoclassical diffusion in tokamaks and stellarators. In particular, each particle with a given transverse invariant in the nonaxisymmetric trap will drift on some drift surface which, in general, will be different for different particles. Thus, unless care is taken with the trap geometry, the drift surfaces of different particles will "split" and intersect. Moreover, if the drift displacement of a particle in the course of one longitudinal oscillation in the trap is large enough, the motion of the particle may in general become stochastic, so that the drift surfaces will also be disrupted. Transverse losses may be reduced by optimizing the magnetic configuration so as to ensure approximate "inscription" (i.e., the absence of crossing) for the drift surfaces. Since, in a nonuniform field, the particles drift in the direction of the binormal, the splitting of the drift surfaces can be entirely avoided by ensuring that the lines of force are geodesics on the magnetic (and therefore drift) surfaces.[123] For long traps, the optimization of the magnetic configuration can be based on the paraxial approximation, and the conditions for no splitting of the drift surfaces can then be formulated.[123]

Ideas concerning an axisymmetric ambipolar trap, in which the problem of enhanced transverse transport does not arise, have recently also been put forward. The attainment of stability is the main problem here. Tubular plasma geometry with "antiplug" field geometry at the ends has been proposed as the solution of this problem.

The use of the ambipolar potential is not the only way of closing the loss cone. One of the alternatives, using the rotation of tubular plasma in a magnetic field of mirror geometry under the influence of a strong radial electric field, has been implemented in the PSP-2 installation (diameter 2 m, magnetic field ~20 kG) at the Institute of Nuclear Physics at Novosibirsk. Stabilization of flute instability is in this case facilitated by nonuniform rotation.

The mirror geometry of the magnetic field is being used not only in traps with magnetic confinement of tenuous plasma, but also in systems with dense plasma (multiplug systems with $\beta \gg 1$, and gas-dynamic systems with a very strong magnetic field in the plugs). In multiplug traps, the magnetic field serves only to reduce thermal conductivity, with the plasma pressure acting directly on the wall.

Finally, hybrid traps incorporating both magnetic and electric fields are being used in plasma confinement. They include, for example, the electromagnetic Yupiter trap at the Khar'kov Physicotechnical Institute. External electrodes are used in these traps to set up separate potential barriers for electrons and for ions and this, in principle, should reduce longitudinal plasma losses. However, there is the attendant danger that this will result in the enhancement of transverse plasma losses due to kinetic instabilities connected with the large equilibrium-pressure drop over a distance shorter than the ion Larmor radius. Experimental work performed on the Atoll toroidal trap at the Kurchatov Institute of Atomic Energy, which has a hyperbolic configuration of magnetic lines of force in its cross-section, has confirmed these fears.

7. CONCLUSIONS

We have been able to give only a brief account of the physical ideas underlying the principal fields of research on magnetic plasma confinement, and of the advances made in this area. A more detailed discussion of the questions touched upon in this paper may be found in recent review articles and monographs.[124-140, 86, 110]

Research on magnetic confinement of high-temperature plasmas has come a long way since those early beginnings during the lifetime of Igor Vasil'evich Kurchatov, when novice researchers were inclined to believe that it would take a few months or, at worst, a few years to take the solution of this problem to a stage where basic difficulties would have been overcome and a clear way could be seen toward the solution of the remaining problems. The fundamentals of the theory of high-temperature plasma and of its confinement by a magnetic field have now been established. The theory is capable of explaining and guiding experimental research. This research now has an extensive and sufficiently well-defined basis, and its own experimental techniques and diagnostic procedures. The production and investigation of plasmas with temperatures of a few keV has become a reality. Installations are being built for experiments capable of demonstrating controlled thermonuclear fusion. At the same time, work has begun on the development of a thermonuclear reactor based on magnetic confinement. We are now seeing the reliable confirmation, and the first steps toward the practical implementation, of the idea of the magnetic thermonuclear reactor that arose more than thirty years ago and was so enthusiatically supported by I. V. Kurchatov who organized research into controlled thermonuclear fusion with the aim of mastering the inexhaustible supply of nuclear energy offered to us by the light elements.

[1] I. E. Tamm, Fizika plazmy i problema upravlyaemykh termoyadernykh reaktsiĭ (Plasma Physics and the Problem of Controlled Thermonuclear Reactions), Izd-vo AN SSSR (USSR Academy of Sciences Press), Moscow, **1**, 3 (1958).

[2] A. D. Sakharov, *ibid.*, **1**, 31 (1958).

[3] I. E. Tamm, *ibid.*, **1**, 31 (1958).

[4] L. A. Artsimovich, A. M. Andrianov, E. I. Dobrokhotov, S. Yu. Luk'yanov, I. M. Podgornyĭ, V. I. Sinitsyn, and I. V. Filippov, Atomnaya Energiya, No. 3, 84 (1956).

[5] I. V. Kurchatov, Uspekhi Fiz. Nauk **59**, 603 (1956).

[6] M. A. Leontovich and V. D. Shafranov, Fizika plazmy i problema upravlyaemykh termoyadernykh reaktsiĭ (Plasma Physics and the Problem of Controlled Thermonuclear Reactions), Izd-vo AN SSSR (USSR Academy of Sciences Press), Moscow, **1**, 207 (1958).

[7] M. Kruskal and J. Tuck, Proc. Roy. Soc. **245A**, 222 (1958).

[8] V. D. Shafranov, Fizika plazmy i problema upravlyaemykh termoyadernykh reaktsiĭ (Plasma Physics and the Problem of Controlled Thermonuclear Reactions), Izd-vo AN SSSR (USSR Academy of Sciences Press), Moscow, **2**, 130 (1958).

[9] V. D. Shafranov, Atomnaya Energiya, No. 5, 38 (1956).

[10] V. D. Shafranov, Zh. Eksp. Teor. Fiz. **33**, 710 (1956) [Sov. Phys. JETP **6**, 545 (1958)].

[11] H. Grad and H. Rubin, in: Proceedings of the Second International Conference on the Peaceful Uses of Atomic Energy, Geneva, 1958 [Russian Translation of selected papers, Atomizdat, Moscow, 1959].

[12] L. I. Artemenkov *et al.*, in: Proceedings of the Sixth European Conference on Controlled Fusion and Plasma Physics, Moscow, **1**, 153 (1973).

[13] A. A. Galeev and R. Z. Sagdeev, Zh. Eksp. Teor. Fiz. **53**, 348 (1967) [Sov. Phys. JETP **26**, 233 (1968)].

[14] A. A. Galeev and R. Z. Sagdeev, in: Voprosy teorii plazmy (Problems in Plasma Theory), ed. M. A. Leontovich, Atomizdat, Moscow, 1973, p. 205.

[15] F. L. Hinton and R. D. Hazeltine, Rev. Mod. Phys. **15**, 239 (1976).

[16] A. A. Ware, Phys. Rev. Lett. **25**, 916 (1970).

[17] R. I. Bickerton, J. W. Connor, and J. B. Taylor, Nature **229**, 110 (1971).

[18] B. B. Kadomtsev and V. D. Shafranov, in: Proceedings of the Fourth International Conference on Plasma Physics and Controlled Nuclear Fusion Research, Madison, 1971, IAEA, Vienna, **II**, 479 (1971).

[19] K. Molvig *et al.*, Comments on Plasma Physics and Controlled Fusion, Comments on Modern Physics **8**, 113 (1982).

[20] L. A. Artsimovich, A. V. Glukhov, and M. P. Petrov, ZhETF Pis. Red. **11**, 449 (1970) [Sov. Phys. JETP Lett. **11**, 304 (1970).

[21] L. A. Artsimovich, ZhETF Pis. Red. **13**, 101 (1971) [Sov. Phys. JETP Lett. **13**, 70 (1971)].

[22] "International tokamak reactor: zero phase", in: Report of the International Tokamak Reactor Workshop, Vienna, 1979, IAEA, Vienna, I, 211 (1980).

[23] B. B. Kadomtsev, Fizika Plazmy, **1**, 531 (1975) [Sov. J. Plasma Phys. **1**, 295 (1979)].

[24]J. W. Connor and J. B. Taylor, Nuclear Fusion 17, 1047 (1977).
[25]V. A. Vlasenkov, V. M. Leonov *et al.*, in: Proceedings of the Seventh International Conference on Plasma Physics and Controlled Nuclear Fusion Research, Innsbruck, 1978, IAEA, Vienna, I, 211 (1979).
[26]V. G. Merezhkin and V. S. Mukhovatov, ZhETF Pis. Red. 33, 463 (1981) [Sov. Phys. JETP Lett. 33, 446 (1981)].
[27]A. B. Berlizov *et al.*, in: Proceedings of the Ninth International Conference on Plasma Physics and Controlled Fusion Research, Baltimore, 1982, CN-41/I-6, IAEA, Vienna, 1983.
[28]F. Alladio *et al.*, *ibid.*, CN-41/I-4.
[29]Alcator Group, *ibid.*, CN-41/I-3.
[30]B. B. Kadomtsev and O. P. Pogutse, in: Proceedings of the Seventh International Conference on Plasma Physics and Nuclear Fusion Research, Innsbruck, 1978, IAEA, Vienna, 1, 649 (1979).
[31]V. V. Parail and O. P. Pogutse, ZhETF Pis. Red. 32, 408 (1980) [Sov. Phys. JETP Lett. 32, 384 (1980)].
[32]T. H. Stix, Nuclear Fusion 18, 353 (1978).
[33]J. D. Callen, Phys. Rev. Lett. 39, 1540 (1977).
[34]A. B. Rechester and M. N. Rosenbluth, Phys. Rev. Lett. 40, 38 (1978).
[35]A. A. Galeev and L. M. Zelenyi, ZhETF Pis. Red. 29, 669 (1979) [Sov. Phys. JETP Lett. 29, 614 (1979)].
[36]V. M. Leonov *et al.*, in: Proceedings of the Eighth International Conference on Plasma Physics and Controlled Fusion Research, Brussels, 1980, IAEA, Vienna, I, 393 (1981).
[37]D. Johnson *et al.*, in: Proceedings of the Ninth International Conference on Plasma Physics and Controlled Fusion Research, Baltimore, 1982, CN-41/A-1, IAEA, Vienna, 1983.
[38]M. Nagumi *et al.*, *ibid.*, CN-41/A-2.
[39]K. Steinmetz *et al.*, *ibid.*, CN-41/A-3.
[40]M. Murakami *et al.*, *ibid.*, CN-41/A-4.
[41]B. B. Kadomtsev, O. P. Pogutse, and E. I. Yurchenko, *ibid.*, CN-41/P-3.
[42]B. A. Carreras, *ibid.*, CN-41/P-4.
[43]A. G. Barsukov *et al.*, in: Proceedings of the Ninth International Conference on Plasma Physics and Controlled Nuclear Fusion Research, Baltimore, 1982, CN-41/A-6, IAEA, Vienna, 1983.
[44]S. Fairfax *et al.*, in: Proceedings of the Eighth International Conference on Plasma Physics and Controlled Nuclear Fusion Research, Brussels, 1980, IAEA, Vienna, I, 450 (1981).
[45]K. H. Burrell *et al.*, Phys. Rev. Lett. 41, 1382 (1978).
[46]R. S. Isler *et al.*, in: Proceedings of the Eighth International Conference on Plasma Physics and Controlled Nuclear Fusion Research, Brussels, 1980, IAEA, Vienna, I, 450 (1981).
[47]TFR Group Report, in: Proceedings of the Ninth International Conference on Plasma Physics and Controlled Nuclear Fusion Research, Baltimore, 1982, CN-41/R-5, IAEA, Vienna, 1983.
[48]H. Maeda *et al.*, in: Proceedings of the Seventh International Conference on Plasma Physics and Controlled Nuclear Fusion Research, Innsbruck, 1978, IAEA, Vienna, I, 377 (1979).
[49]A. V. Bortnikov *et al.*, Proceedings of the Eighth International Conference on Plasma Physics and Controlled Nuclear Fusion Research, Brussels, 1980, IAEA, Vienna, I, 687 (1981).
[50]S. K. Erents *et al.*, *ibid.*, I, 697 (1981).
[51]M. Nagami *et al.*, *ibid.*, II, 367 (1981).
[52]M. Keilhacker *et al.*, *ibid.*, 351 (1981).
[53]V. D. Shafranov, Zh. Tekh. Fiz. 40, 241 (1970) [Sov. Phys. Tech. Phys. 15, 175 (1970)].
[54]L. E. Zakharov, Fizika Plazmy, 6, 18 (1975) [*sic*].
[55]Yu. N. Dnestrovskiĭ, D. P. Kostomarov, G. V. Pereverzev, and K. N. Tarasyan, Fizika Plazmy, 4, 1001 (1978) [Sov. J. Plasma Phys. 4, 557 (1978)].
[56]B. B. Kadomtsev, Fizika Plazmy, 1, 710 (1975) [Sov. J. Plasma Phys. 1, 389 (75)].
[57]V. D. Shafranov and E. I. Yurchenko, Zh. Eksp. Teor. Fiz. 53, 1157 (1967) [Sov. Phys. JETP 26, 682 (1968)].
[58]J. W. Connor, R. J. Hastie, and J. B. Taylor, Phys. Rev. Lett. 40, 396 (1978).
[59]O. N. Babishchevich, L. M. Degtyarev, and S. Yu. Medvedev, Preprint IMP AN SSSR, No. 154; Fizika Plazmy 9, No. 1 (1983) [Sov. J. Plasma Phys. (in press)].
[60]A. M. Todd, J. Manicam, M. Akabayashi, *et al.* Nucl. Fusion 19, 743 (1979).
[61]S. von Goeler, W. Stodiek, and N. Sauthoff, Phys. Rev. Lett. 33, 1201 (1974).
[62]B. V. Waddel, D. A. Monticello, M. N. Rosenbluth, and R. B. White, Nucl. Fusion 16, 528 (1976).
[63]A. F. Danilov, Yu. N. Dnestrovskiĭ, D. P. Kostomarov, and A. M. Popov, Fizika Plazmy 2, 167 (1976) [Sov. J. Plasma Phys. 2, 93 (1976)].
[64]B. A. Carreras, H. R. Hicks, J. A. Holmes, B. V. Waddel, Phys. Fluids 23, 1811 (1980).
[65]B. B. Kadomtsev and O. P. Pogutse, Zh. Eksp. Teor. Fiz. 6 65, 575 (1973) [Sov. Phys. JETP 38, 283 (1974)].
[66]D. W. Ross and S. M. Mahayan, Phys. Rev. Lett. 40, 324 (1978).
[67]K. Tsang *et al.*, *ibid.*, 40, 327 (1978).
[68]L. Chen, *ibid.*, 41, 649 (1978).
[69]T. M. Antonsen jnr, *ibid.*, 41, 33 (1978).
[70]B. B. Kadomtsev and O. P. Pogutse, Nucl. Fusion 10, 67 (1971).
[71]B. B. Kadomtsev and O. P. Pogutse, Zh. Eksp. Teor. Fiz. 53, 2025 (1967) [Sov. Phys. JETP 26, 1146 (1968)].
[72]V. V. Parail and O. P. Pogutse, Fizika Plazmy 2, 228 (1976) [Sov. J. Plasma Phys. 2, 125 (1976)].
[73]V. V. Parail and O. P. Pogutse, Nucl. Fusion 18, 303 (1978).
[74]V. V. Alikaev, Yu. I. Arsen'ev, G. A. Bobrovskiĭ, *et al.*, Zh. Tekh. Fiz. 45, 515 (1975) [Sov. Phys. Tech. Phys. 20, 322 (1975)].
[75]V. V. Alikaev, K. A. Razumova, Yu. A. Sokolov, Fizika Plazmy 1, 546 (1971) [Sov. J. Plasma Phys. 1, 303 (1971)].
[76]C. Bernabei *et al.*, in: Proceedings of the Ninth International Conference on Plasma Physics and Controlled Nuclear Fusion, Baltimore, 1982, CN-41/C-5, IAEA, Vienna, 1983.
[77]V. V. Alikaev, in: Proceedings of the Tenth European Conference on Controlled Fusion and Plasma Physics, Moscow, 2, 11 (1953).
[78]R. A. Demirkhanov, A. G. Kirov, L. F. Ruchko, and A. V. Sukhanov, *ibid.*, 1, paper E-7.
[79]V. V. Alikaev, V. L. Vdovin, D. P. Ivanov, *et al.*, in: Proceedings of the Ninth International Conference on Plasma Physics and Controlled Nuclear Fusion, Baltimore, 1982, CN-41/K-1, IAEA Vienna, 1983.
[80]M. Porcolab *et al.*, *ibid.*, CN-41/C-4.
[81]L. Spitzer Jr., Phys. Fluids 1, 253 (1958).
[82]J. M. Greene, J. L. Johnson, and K. E. Weimer, Phys. Fluids 8, 145 (1962).
[83]V. D. Pustovitov, Fizika Plazmy 8, 473 (1982) [Sov. J. Plasma Phys. 8, 265 (1982)].
[84]L. E. Zakharov, V. D. Pustovitov, V. D. Shafranov, L. M. Degtyarev, *et al.*, in: Proceedings of the Ninth International Conference on Plasma Physics and Controlled Nuclear Fusion, Baltimore, 1982, CN-41/V-5, IAEA, Vienna, 1983.
[85]M. N. Rosenbluth, R. Z. Sagdeev, J. B. Taylor, G. M. Zaslavskiĭ, Nucl. Fusion 6, 297 (1966).
[86]M. S. Rabinovich, in: Itogi nauki i tekhniki (Advances in Science and Technology), ser. Fizika Plazmy, Moscow, VINITI, 2, 6 (1981).
[87]A. G. Dikiy *et al.*, in: Proceedings of the Sixth European Conference on Controlled Fusion and Plasma Physics, Moscow, 1, 105 (1973).
[88]G. Cattanei, in: Proceedings of the Ninth International Conference on Plasma Physics and Controlled Nuclear Fusion Research, Baltimore, 1982, CN-41/L-5, IAEA, Vienna, 1983.
[89]K. Uo *et al.*, *ibid.*, CN-41/L-3.
[90]L. M. Kovrizhnykh, Preprints FIAN SSSR, Nos. 165 and 222, Moscow, 1982.

[91]R. E. Potok, P. A. Politzer, and L. M. Lidsky, Phys. Rev. Lett. **45**, 1328 (1980).
[92]A. N. Boozer and G. Kuo-Petravic, Phys. Fluids **24**, 851 (1981).
[93]H. E. Mynick, Phys. Fluids **25**, 325 (1982).
[94]H. E. Mynick, T. K. Chu, and A. H. Boozer, Phys. Rev. Lett. **48**, 322 (1982).
[95]S. N. Popov, and A. P. Popryadukhin, Zh. Tekh. Fiz. **36**, 390 (1966) [Sov. Phys. Tech. Phys. **11**, 284 (1966)].
[96]M. A. Ivanovskiĭ, S. N. Popov, A. P. Popryadukhin, and M. S. Rabinovich, in: Proceedings of the Fourth International Conference on Plasma Physics and Controlled Nuclear Fusion Research, Madison, 1971, IAEA, Vienna, **3**, 63 (1972).
[97]J. F. Lyon *et al*., in: Proceedings of the Ninth International Conference on Plasma Physics and Controlled Nuclear Fusion Research, Baltimore, 1982, CN-41/Q-3, IAEA, Vienna, 1983.
[98]U. Brossman *et al*., *ibid*., CN-41/Q-5.
[99]V. V. Arsenin *et al*., *ibid*., CN-41/Q-6.
[100]A. H. Boozer *et al*., *ibid*., CN-41/Q-4.
[101]P. C. Thonemann *et al*., Nature, **181**, 217 (1958).
[102]A. P. Babichev *et al*., Nucl. Fusion, **2**, 84 (1962); Nucl. Fusion Supplement **2**, 635 (1962).
[103]B. B. Kadomtsev, *ibid*., **3**, 969 (1962).
[104]D. C. Robinson, Plasma Phys. **13**, 439 (1971).
[105]J. B. Taylor, in: Pulsed High-Beta Plasmas, Pergamon Press, Oxford, 1976, p. 59.
[106]T. Ohkava, M. Chu, and M. Schaffer, Nucl. Fusion **20**, 1464 (1980).
[107]D. A. Baker *et al*., in: Proceedings of the Ninth International Conference on Plasma Physics and Controlled Nuclear Fusion, Baltimore, 1982, CN-41/H-2-1, IAEA, Vienna, 1983.
[108]V. V. Belikov *et al*., *ibid*., CN-41/M-2-1.
[109]W. T. Armstrong *et al*., *ibid*., CN-41/M-2-1.
[110]V. A. Burtsev, V. A. Gribkov, and T. I. Filippova, in: Itogi nauki i tekhniki (Advances in Science and Technology), ser. Fizika Plasmy, VINITI, Moscow, **2**, 80, 1981.
[111]Yu. V. Gott, M. S. Ioffe, and V. G. Tel'kovskii, Nucl. Fusion Supplement 3, 1045 (1962).
[112]Yu. T. Baiborodov *et al*., in: Proceedings of the Sixth European Conference on Controlled Fusion and Plasma Physics, Moscow, **2**, 122 (1973).
[113]M. S. Ioffe, B. I. Kanaev, V. P. Pastukhov, and E. E. Yushmanov, Zh. Eksp. Teor. Fiz. **67**, 2145 (1974) [Sov. Phys. JETP **40**, 1064 (1975)].
[114]F. H. Coensgen *et al*., Phys. Rev. Lett. **35**, 1501 (1975).
[115]F. H. Coensgen *et al*., *ibid*., **37**, 143 (1976).
[116]G. I. Dimov, V. V. Zkakaidakov, and M. E. Kishinevskiĭ, Fizika Plazmy **2**, 597 (1976) [Sov. J. Plasma Phys. **2**, 326 (1976)].
[117]T. K. Fowler and B. G. Logan, Comments on Plasma Physics and Controlled Fusion, **2**, 167 (1977).
[118]T. C. Simonen *et al*., in: Proceedings of the Ninth International Conference on Plasma Physics and Controlled Fusion Research, Baltimore, 1982, CN-41/G-1, IAEA, Vienna, 1983.
[119]D. E. Baldwin and B. G. Logan, Phys. Rev. Lett. **43**, 1318 (1979).
[120]D. D. Ryutov and G. V. Stupakov, ZhETF Pis. Red. **26**, 186 (1977) [JETP Lett. **26**, 174 (1977)].
[121]D. D. Ryutov and G. V. Stupakov, Fizika Plazmy **4**, 501 (1978) [Sov. J. Plasma Phys. **4**, 278 (1978)].
[122]D. A. Panov, Preprint IAE-3535/6, Moscow, 1982; Fizika Plazmy **9**, 1 (1983) [Sov. J. Plasma Phys. (in press)].
[123]G. V. Stupakov, Fizika Plazmy, **5**, 958 (1979) [Sov. J. Plasma Phys. **5**, 534 (1979)].
[124]V. S. Mukhovatov, in: Itogi nauki i tekhniki (Advances in Science and Technology), ser. Fizika Plazmy, VINITI, Moscow, 1980, **1**, p. 6.
[125]V. A. Chuyanov, *ibid*., p. 119.
[126]J. Sheffield, Proc. IEEE, **69**, 867 (1981).
[127]T. C. Simonen, *ibid*., **69**, 935 (1981).
[128]G. Beitman, MGD neustoichivosti (MHD Instabilities), Energoizdat, Moscow, 1982.
[129]D. D. Ryutov, Nucl. Fusion **20**, 1068 (1980).
[130]R. J. Bickerton, *ibid*., **20**, 1072 (1980).
[131]V. D. Shafranov, *ibid*., **20**, 1075 (1980).
[132]P. H. Rutherford, *ibid*., **20**, 1086 (1980).
[133]M. N. Rosenbluth and P. H. Rutherford, in: Fusion, ed. E. Teller, Academic Press, New York, 1981, Part 1A, p. 32.
[134]H. P. Furth, *ibid*., p. 124.
[135]J. L. Shohet, *ibid*., p. 243.
[136]T. K. Fowler, *ibid*., p. 291.
[137]R. F. Post, *ibid*., p. 358.
[138]D. A. Baker and W. E. Quinn, *ibid*., p. 438.
[139]H. A. B. Bodin, in: Proceedings of the Tenth International Conference on Controlled Fusion and Plasma Physics, Moscow, **2**, p. 21 (1981).
[140]V. D. Shafranov, Phys. Fluids, **26**, No. 2 (1983).

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