Open-ended traps

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The open-ended trap is one of the installations used for the magnetic confinement of thermonuclear plasma. Open-ended traps have a number of important advantages as compared with other confinement systems: they are attractive from the engineering point of view, their magnetic field is efficiently used to confine the plasma, they can be operated under steady-state conditions, and there is no particular problem with removing thermonuclear reaction products and heavy impurities from the plasma. At the same time, it has long been considered that the open-ended trap has a doubtful future as a basis for a thermonuclear reactor because of the relatively high rate of loss of plasma along the magnetic lines of force. The situation has changed for the better during the last decade, and a number of improved traps, that are largely free from this defect, has been proposed. This review examines the physical principles of open-ended traps (ambipolar, centrifugal, multimirror, gas-dynamic, and so on), the present state of research into these systems, and their future prospects. The use of open-ended traps as high-flux generators of 14-MeV neutrons is also discussed.

CONTENTS

1. INTRODUCTION

For a steady-state thermonuclear fusion reaction to occur in plasma at 100 million degrees, the plasma must be thermally insulated from the objects surrounding it. This is achieved in the so-called magnetic traps which are commonly divided into two major classes, namely, open-ended and closed traps. In closed traps, the magnetic lines of force do not cross the plasma boundaries and the plasma is confined in a torus (Fig. 1a). In open-ended traps, on the other hand, the plasma confinement region is bounded in the direction of the lines of force, and usually looks like a truncated cylinder, deformed at the ends. The best-known example of an open-ended trap is the mirror trap (known as the *probkotron* (probka = cork) in Russian literature; see Fig. 1b). The principle of this trap was first put forward in the 1950s by G. I. Budker in the USSR[1] and, independently, by R. F. Post in the USA.[2] There are many other varieties of the open-ended trap, but they all include the basic elements of the mirror trap to a greater or lesser extent.

Research into open-ended traps has not had a smooth history. Early calculations by G. I. Budker had already demonstrated that the rate of loss of plasma along a magnetic line of force was relatively high, so tnat, from the purely thermonuclear point of view, the future of the mirror trap did not look very bright. It seemed that interest in these devices would decline. In actual fact, the mirror trap became one of the most popular plasma installations in the 1950s and early 1960s. The point was that thermonuclear power seemed (and was!) a very remote possibility, and did not dictate, to the extent that it does now, the choice of the system to be investigated. Technological simplicity and experimental flexibility, and also the provision of facilities for an extensive range of physical investigations, were more important at the time and, in this respect, the open-ended trap is difficult to compete with.

It is precisely the experiments with open-ended traps, performed in 1960s, that have contributed the major part of what is now regarded as the foundation of plasma physics. Here we mention only the experiments of Ioffe *et al.*[3] who were the first to demonstrate the feasibility of MHD stabilization of plasmas by the "magnetic well" method, the work of Zavoiskiĭ *et al.*[4] who investigated the so-called turbulent plasma heating, and the stabilization of large-scale plasma instabilities by the feedback method at Golovin's Laboratory.[5]

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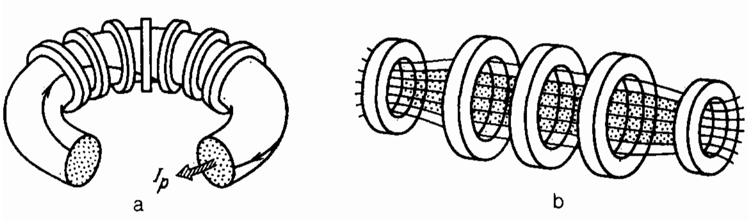

FIG. 1. Comparison of closed and open-ended traps: a—closed tokamak-type trap; the diagram shows part of the plasma ring surrounded by coils producing the toroidal magnetic field; thick arrow shows the direction of the toroidal plasma current whereas the thin arrow shows one of the magnetic lines of force; the lines of force are "wound" on the toroidal surfaces and remain inside the vacuum chamber; b—open-ended trap; the diagram shows coils, magnetic lines of force, and confined plasma.

This was the golden age of the open-ended trap: the search was geographically extensive, and the participation of leading laboratories assured a steady flow of results. Suffice is to say that at the thermonuclear conferences held by the International Atomic Energy Agency (the major forum for the plasma community), open-ended traps accounted for about one third of the proceedings, whereas tokamaks were dealt with in about one-tenth of the contributions.

However, by the late 1960s, the steady advance of the tokamak forced a review of all types of thermonuclear devices as possible candidates for a prototype thermonuclear reactor. It was at this point that the "congenital defect" of the open-ended trap, i.e., its relatively high rate of loss of plasma in the direction of the field, began to show itself. Careful calculations revealed that the energy multiplication factor Q, i.e., the ratio of the generated thermonuclear power to the power expended in heating the plasma, could hardly be made more than 1.2–1.5 under the most optimistic assumptions. This was clearly inadequate for an economically attractive thermonuclear reactor. Attempts to save the situation through the development of special recuperators of the energy transported into the mirrors were not successful because they led to larger and more costly installations.

Although the traditional advantages of the open-ended trap (simple geometry, high plasma pressure in moderate magnetic field, steady-state operation, etc.) were as obvious as before, interest in such traps began to decline. The geographical base of this research contracted,[1)] and there was a decline in the number of physicists working on open-ended traps.

It seemed that this line of work would soon meet the same fate as the previously very respectable studies of pulsed theta pinches, which ceased in the mid-1970s. Only three or four laboratories in the USSR and the USA continued to maintain an interest in such studies. Although the time was not opportune for this, there were constant advances in the experimental technology necessary for open-ended traps, such as injectors of fast hydrogen atoms, used to introduce matter and energy across the magnetic field, complex non-axially symmetric magnetic systems for the MHD stabilization of plasmas, vacuum technology capable of pumping out large quantities of gas, and many other systems. At the same time, attempts were continuing to reduce the longitudinal loss of plasma. The many suggestions included a multimirror system for plasma confinement[6,7] in which the single trap was replaced by a set of coupled traps whose overall length was greater than the mean free path of the ions, a trap with centrifugal plasma confinement[8], a field-reversed system (see Ref. 9)[2)], and certain other systems. Unfortunately, the first of these required technological facilities that were regarded as too exotic in the early 1970s, whereas the second and third were based on assumptions about MHD plasma stability that were not immediately obvious.

A turn for the better in the study of open-ended traps was signalled by two important events that occurred in the mid 1970s.

First, the available experimental techniques were clearly shown in 1975-1976 to be effective in experiments on the 2XIIB installation at the Livermore Laboratory in the USA, in which extensive use was made of injectors of neutrals, plasma jet techniques, and titanium gettering of all surfaces facing the plasma. Quasistationary plasmas with densities of 10^{14} cm^{-3} and ion "temperature" of 10–12 keV (Ref. 10) were produced by injecting a deuterium atomic beam of 7 MW power.[3)]

Second, in the middle of 1976, Dimov, Zakaĭdakov, and Kishenevskiĭ,[11] and, in the beginning of 1977, Fowler and Logan,[12] independently published the principle of the so-called ambipolar trap which reduced the rate of longitudinal plasma loss when it was used together with the techniques that were successfully employed on the 2XIIB system.

The fact that these two events occurred at more or less the same time led to renewed interest in open-ended traps, increased financing for their development, and a new expansion of their geographical base. Major ambipolar traps began to be built in the USA, USSR, and Japan. It was soon demonstrated that ambipolar confinement did work and produced a significant reduction in the loss of plasma through the mirrors. Still larger versions of the open-ended trap were then built and planned.

However, certain specific difficulties emerged in the course of time. First, the large-scale magnetohydrodynamic stability of plasmas is assured in all the installations built so far by the use of the non-axially-symmetric magnetic fields that had previously proved themselves so well in the past. It emerged,[13] however, that they can easily give rise to enhanced transverse plasma loss, which becomes particularly appreciable when the longitudinal loss is suppressed. Second, not all was well with longitudinal confinement either: for reasons that are still not totally understood, attempts to increase the plasma density resulted in a deterioration in this type of confinement. There is a suspicion that this was due to development of plasma microfluctuations. Thus, on the one hand, intensive searches are now in progress for axially symmetric MHD-stable plasma configurations, and, on the other hand, attempts are being made to understand the phenomena that determine the growth of longitudinal loss at high plasma density.

Alongside the complications that have arisen in connection with ambipolar traps, the overall picture has some important positive elements, e.g., the discovery of the considerable latent potentialities of other open-ended trap systems. In particular, it has been found that an increase in the length and mirror ratio of a two-mirror trap, produces a substantial increase in the longitudinal plasma lifetime. A system based on this approach has been referred to as the gas-dynamic trap (GDT).[14] Moreover, it has become clear that there are several ways in which the length of multimirror traps can be reduced and their parameters can be brought significantly closer to the possibilities offered by modern technology. Finally, systems with rotating plasmas have been found to have further plasma-stabilizing potentialities. It is important to note at this point that these lines of research into open-ended traps are still at the stage at which initial experimental data are being acquired, and it is still too early to say much about their role in the future.

It has also become clear in recent years that open-ended traps can be used as neutron generators for tests in materials science. The designs that are being developed for such generators aim for plasma parameters that are not too different from those already achieved experimentally. The development of such devices may be looked upon as a useful intermediate problem.

On the whole, research into open-ended traps is now at a very decisive stage. The next year or two will show whether the hopes invested in such systems are justified, and whether this approach to the problem of controlled thermonuclear fusion will have to await better times, by which experimental techniques and available technologies will have advanced, and a better understanding of the properties of plasmas will have been achieved.

This review presents a brief description of the basic types of open-ended traps and of the results obtained with them. The bibliography given at the end is in no way exhaustive. The reader will find further information in the review literature.[15–19]

2. MIRROR TRAPS

2.1. Rate of loss of plasma through the mirrors

The reflection of charged particles by a magnetic mirror is due to the adiabatic invariance of the quantity

$$\mu=\frac{mv_\perp^2}{2B}, \tag{1}$$

where m is the particle mass, $v_\perp$ is the transverse (relative to the magnetic field) velocity component, and B is the magnetic field. The invariant μ can be interpreted as the magnetic moment of the "Larmor circle." The degree to which μ is conserved is found to increase as the Larmor radius of the particle (calculated from its total energy) is reduced as compared with the characteristic scale of changes in the magnetic field, i.e., $L\equiv B/|\nabla B|$. In practice, when the Larmor radius is less than L by a factor of 5–6, μ can be looked upon in the present context simply as a constant of motion. We shall actually make this assumption. A detailed examination of the conservation of μ can be found, for example, in the review by B. V. Chirikov.[20]

Using the law of conservation of energy

$$\mathscr{E}=\frac{m(v_\parallel^2+v_\perp^2)}{2B}=\text{const}, \tag{2}$$

and the expression given by (1), we can show that the variation in the longitudinal velocity of a particle during its motion along a line of force is described by

$$\frac{mv_\parallel^2}{2}=\mathscr{E}-\mu B(s)\equiv\mathscr{E}-U(s), \tag{3}$$

where s is the coordinate measured along the line of force. The function $U(s)$ is the effective potential energy associated with longitudinal motion. Since $\mathscr{E}$ and μ do not depend on time, the condition that the particle will be confined within the trap is obviously that the maximum of $U(s)$ must be greater than $\mathscr{E}$, i.e.,

$$\mathscr{E}-\mu B_{\max}<0 \tag{4}$$

where $B_{\max}$ is the maximum magnetic field in the mirror.

Using the subscript 0 to denote all quantities at the center of the magnetic mirror trap (in the uniform part of the magnetic field), we can rewrite (4) in the form

$$\frac{v_{\perp0}^2}{v_{\perp0}^2+v_{\parallel0}^2}>\frac{B_0}{B_{\max}}. \tag{5}$$

The quantity

$$R\equiv\frac{B_{\max}}{B_0}$$

is commonly referred to as the mirror ratio. Let θ_0 be the angle between the particle velocity vector and the magnetic field within the uniform field region. We can then rewrite (5) in the form

$$\theta_0>\arcsin R^{-1/2}. \tag{6}$$

The last two inequalities show that the velocity-space region from which particles are lost is a cone with axis parallel to the magnetic field (Fig. 2). This is referred to as the loss cone. Particles lying outside the loss cone oscillate between the mirrors (with $\mu=$ const satisfied sufficiently accurately), and are confined to the trap for practically an indefinite time. The particles leave the trap only as a result of scattering by one another, which produces a change in the angle θ_0, and, in the final analysis, the entry of the velocity vector into the loss cone.

Let us estimate the rate at which the ions are lost. We know that Coulomb scattering of plasma particles has the Fokker-Planck (diffusion) character. Accordingly, the time necessary for the scattering of charged particles through an angle that is not too large, i.e., $\theta\lesssim1$, is proportional to θ^2. For ions, this time can be estimated from the formula (see Ref. 21)

$$\tau_{ii}^{(\theta)}\sim\frac{\theta^2W_i^2}{\pi\Lambda ne^4}\left(\frac{m_i}{2W_i}\right)^{1/2}, \tag{7}$$

where n is the number of ions per unit plasma volume, W_i is their characteristic energy, e is the charge of the electron, and Λ is the so-called Coulomb logarithm which, under typical thermonuclear conditions, is approximately equal to 15.

Substituting $\theta=1$ in (7), we obtain the ion-ion collision time

$$\tau_{ii}\equiv\frac{W_i^2}{\pi\Lambda ne^4}\left(\frac{m_i}{2W_i}\right)^{1/2}. \tag{7'}$$

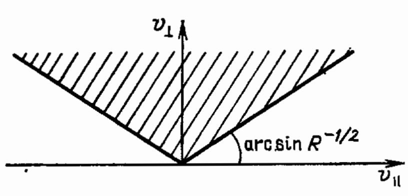

FIG. 2. Confinement region in velocity space (shaded).

The product of τ_{ii} and the ion velocity gives the ion mean free path:

$$\lambda_{ii} \sim \frac{W_i^2}{\pi \Lambda n e^4}.$$

In the hypothetical thermonuclear reactor in the form of a mirror trap, λ_{ii} is much greater than the length of the system (for $W_i \sim 100$ keV and $n \sim 10^{14}$ cm^{-3}, we have $\lambda_{ii} \sim 3 \cdot 10^5$ m).

For mirror ratios not too close to unity, $R - 1 \sim 1$, the angle of the loss cone is of the order of unity, i.e., [see (7)]

$$\tau \sim \tau_{ii}. \tag{8}$$

For the same energy as that of the ions, the electron-ion and electron-electron scattering times τ_{ei} and τ_{ee} are smaller than the ion scattering time by a factor of about $(m_i/m_e)^{1/2}$ [see Ref. 21; this can also be seen from (7′) if we replace m_i with m_e]. Accordingly, the electrons rapidly become isotropic (and, moreover, their distribution function becomes Maxwellian), and they fill the loss cone. However, since the plasma must remain quasineutral, the initial more rapid loss of electrons produces a certain positive potential in the plasma, and the spatial distribution of this potential adjusts itself so that the quasineutrality condition is satisfied locally throughout the plasma.[4)] The potential that evolves in this way is called the ambipolar potential.

Since, as already noted, the electron distribution function is Maxwellian, the ambipolar potential can also be expressed in terms of the electron density by the Boltzmann formula

$$e\varphi(s) = T_e \ln n(s) + \text{const}, \tag{9}$$

where T_e is the electron temperature (constant along the line of force because of the high electron thermal conductivity). The plasma density decreases from the center of the trap toward the mirrors. Accordingly, the ambipolar electric field is given by

$$E_{\parallel} = -\frac{\partial \varphi}{\partial s} = -\frac{T_e}{n}\frac{\partial n}{\partial s}$$

and acts on the ions in the direction of the mirrors, producing a deterioration in ion confinement.

When the ambipolar potential is present, the longitudinal motion of the ions is determined not by (3) but by

$$\frac{m_i v_{\parallel}^2}{2} = \mathscr{E} - U_{\mathrm{Yu}}(s), \tag{10}$$

where

$$\mathscr{E} = \frac{m_i v^2}{2} + e\varphi = \text{const}$$

is the total energy of an ion and

$$U_{\mathrm{Yu}}(s) = e\varphi(s) + \mu B(s) \tag{11}$$

is the effective potential energy associated with longitudinal motion, called the Yushmanov potential. The problem of whether or not the particle is confined within the trap depends on the relationship between $\mathscr{E}$ and the maximum of $U_{\mathrm{Yu}}(s)$. Figure 3a shows typical graphs of $\varphi(s)$ and $B(s)$. It is clear from the figure and from (11) that the point at which U_{Yu} is a maximum is not in general the same as the point at

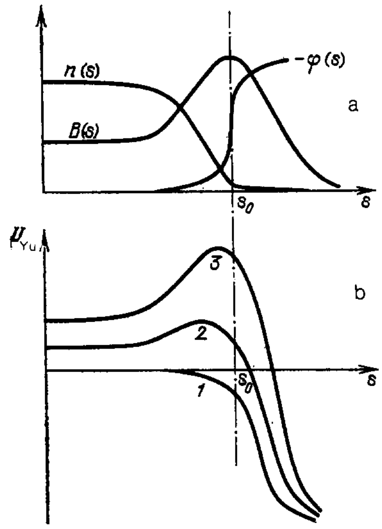

FIG. 3. The effect of an electric field on particle confinement in a magnetic mirror trap: a—general shape of the functions $B(s)$, $n(s)$, and $\varphi(s)$, assuming that the potential is zero at the center of the trap; the point s_0 corresponds to the maximum of the magnetic field; b—general form of $U_{\mathrm{Yu}}(s)$ for a few values of μ; the curves are numbered in order of increasing μ ($\mu = 0$ for curve 1). The figures are symmetric with respect to the $s = 0$ plane.

which the magnetic field in the mirror is a maximum. Hence to determine the shape of the confinement region in the phase space of the ions, we must know the function $\varphi(s)$ or, equivalently, the function $n(s)$. On the other hand, before we can determine the function $n(s)$ we must know the ion distribution function for which, in turn, we must know the shape of the confinement region. This gives rise to a complicated self-consistency problem that can be solved only by numerical techniques.

It is qualitatively clear that the ambipolar electric field, which extracts ions from the trap, reduces their lifetime. On the other hand, if the potential difference between the trap center and the mirror does not exceed W_i/e [which occurs when the electron temperature is not too high; see (9)], the estimate given by (8) remains valid to within an order of magnitude. Simplified calculations are occasionally based on a model in which the magnetic field and potential distributions along the trap axis are replaced by step functions (Fig. 4), where the jump in the potential occurs on the inner side of the mirror. If we represent the potential drop between the mirror and the center by $\Delta\varphi(\Delta\varphi > 0)$, the longitudinal velocity component of an ion after it crosses this potential drop (but is still inside the mirror) is $[v_{\parallel 0}^2 + (2e\Delta\varphi/m_i)]^{1/2}$, and the transverse component is $v_{\perp 0}$. Accordingly,

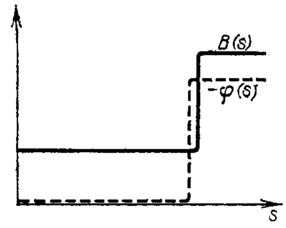

FIG. 4. Step model of the magnetic field and the potential.

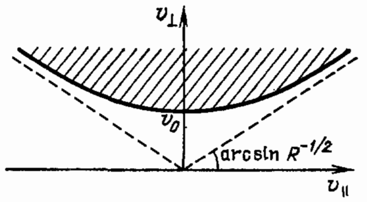

FIG. 5. Confinement region with the ambipolar electric field taken into account for the "step" model of Fig. 4 [$v_0 = (2e\Delta\varphi/Rm_i)^{1/2}$].

instead of (5), we now have the following confinement condition:

$$v_{\perp 0}^2\left(1-\frac{1}{R}\right)-\frac{v_{\parallel 0}^2}{R}>\frac{2e\Delta\varphi}{m_1 R}. \qquad (12)$$

The boundary of the confinement region (12) is a hyperboloid (Fig. 5). Ions with energies below a certain minimum value (equal to $e\Delta\varphi/R$) are not confined within the trap (they are extracted by the electric field). It follows that, in any case, the initial energy of ions introduced into the trap must exceed $e\Delta\varphi/R$:

$$W_1>\frac{e\Delta\varphi}{R}. \qquad (13)$$

The plasma density in a mirror is much lower than at the center of the trap. It follows that, according to (9), the quantity $\Delta\varphi$ amounts to a few electron temperatures (φ depends logarithmically on n). The negative effect of the ambipolar potential can be minimized, as required, by increasing the mirror ratio. If the estimate given by (13) is satisfied to within a considerable margin, the ambipolar field will have little effect on plasma confinement.

The estimate given by (8) refers to the case where $R-1\sim 1$. It must be modified for large mirror ratios. The solution of the transport equation for ions shows that, when $R\gg 1$, the lifetime increases with R only logarithmically

$$\tau\approx 0.4\tau_{ii}\ln R. \qquad (14)$$

Consequently, a significant increase in the lifetime as compared with (8) cannot be achieved by increasing the mirror ratio. The weak dependence of the lifetime on the mirror ratio is related to the Fokker-Planck character of the Coulomb collisions.

The interesting (and useful for thermonuclear applications) property of the magnetic mirror trap is that heavy impurities are poorly confined, or not confined at all. This follows from (13): for an ion with charge Z, the right-hand side of this expression must be replaced with $Ze\Delta\varphi/R$, i.e., by an amount exceeding the "thermal" energy of the ions for sufficiently high values of Z. Such ions will leave the trap in one transit between the mirrors, and their concentration in the plasma will be vanishingly small. For ions with low values of Z, the ambipolar potential does not lead to an immediate escape through the mirrors, but their concentration is also low because of the expansion of the loss hyperboloid, and the higher frequency of scattering of such ions by the main plasma ions. On the whole, it may be concluded that, in the steady state, the plasma in the mirror trap will be very pure. In this respect, the mirror trap is different from closed systems in which the accumulation of impurities in the plasma may be a source of serious difficulty.

2.2. Steady state of plasma in a mirror trap

The loss of particles through the mirrors can be compensated by constantly feeding the trap with new particles. This is usually done by injecting beams of atomic hydrogen (or a hydrogen isotope).

The atomic-beam production technique has advanced rapidly during the last two decades, and deserves more detailed explanation at this point. An atomic beam is produced in three stages. The first step is to prepare slow hydrogen ions in a gas-discharge ion source. These ions are then accelerated to the required energy and are passed through a charge-transfer target which is usually in the form of a cloud of gaseous hydrogen. This neutralizes the fast hydrogen ions via the transfer reaction

$$H^{+*}+H_2\rightarrow H_2^{*}+H^{*} \qquad (15)$$

where the asterisk labels the fast particles. Because the electron mass is low, the momentum of the hydrogen atom is almost equal to that of the original proton. This means that, when the angular divergence of the original proton beam is small, the divergence of the resulting atomic hydrogen beam is also small. Sources of atomic deuterium that are capable of producing deuteron beams with energies in excess of 100 keV have now been developed. The equivalent current is 50 A for a pulse length of several seconds. [22] By using several such sources, the power injected into large tokamaks has been increased to 20 MW. The next step will be to increase the energy and the pulse length still further (continuous beams are the ultimate aim).

At deuteron energies appreciably in excess of 100 keV, the charge-transfer cross section decreases rapidly, and the reaction (15) becomes ineffective. High-energy atomic beams are produced from negative hydrogen ions. These ions are extracted from a special ion source and their negative charges are stripped off in a gas target. This method is now being used to produce beams of hydrogen atoms with energies of several hundred keV, which is not at all the ultimate limit.

The advantage of atomic beams is that they can be introduced into the plasma across a strong magnetic field. The capture of a fast atom by plasma relies on its ionization by electrons and ions, or on charge transfer to plasma ions. [5)] Experiments in which the steady state of the plasma is maintained by the injection of an atomic beam must be arranged so that the capture length for neutrals must be comparable with the thickness of the plasma: if the former is significantly greater than the latter, the beam will pass freely through the plasma, and will be lost. On the other hand, when the reverse condition is satisfied, the beam will not succeed in penetrating the inner plasma layers, and the plasma will become tubular.

The mean energy of plasma ions in this method of maintaining the steady state is a certain fraction (of the order of unity) of the injected energy W_{inj}. The electron temperature is determined by two factors, namely, the transfer of energy from ions to electrons, and the loss of energy through the mirrors. When the plasma is well confined in the trap, its density both within the mirror and outside the mirror is much lower than within the trap and, in accordance with (9), electrons leaving the trap must overcome a high potential barrier of height equal to a few electron temperatures. Accordingly, the loss of each electron from the system is accompanied by a loss of energy equal to $A_e T_e$, where the coefficient A_e is typically 6–8.

The power transmitted from ions to electrons per unit plasma volume is (see Ref. 21)

$$n\,\frac{\overline{W}_i-(3/2)\,T_e}{\tau_{ie}^{E}}, \tag{16}$$

where

$$\tau_{ie}^{E}=\frac{3m_iT_e^{3/2}}{8m_e^{3/2}\sqrt{2\pi}\Lambda e^4 n} \tag{17}$$

can be interpreted as the time necessary for energy transfer between ions and electrons. Since, in accordance with the quasineutrality condition, the electron and ion lifetimes are equal, the electron energy loss per unit plasma volume is $A_e T_e n/\tau$, where τ is the ion lifetime that can be estimated from (14). Accordingly, in equilibrium, we have

$$\frac{\overline{W}_i-(3/2)\,T_e}{\tau_{ie}^{E}}=\frac{A_eT_e}{\tau}.$$

Hence, it readily follows that

$$\frac{T_e}{\overline{W}_i}\sim\left(\frac{m_e}{m_i}\right)^{1/5}\left(\frac{\ln R}{A_e}\right)^{2/5}, \tag{18}$$

i.e., the electron temperature is significantly lower than the mean ion energy. The temperature T_e is lower by an order of magnitude than $\overline{W}_i$. At the same time, the time τ_{ie}^{E} for the slowing down of ions by electrons is comparable with τ, so that quantitative calculations on the confinement of ions within a trap must take the drag by electrons into account.

Occasionally, the stability of plasma with respect to the excitation of microfluctuations (see Section 2.4 for further details) is maintained by passing a beam of relatively cold plasma through the trap. This cold-plasma beam is prepared outside the trap in special gas-discharge sources. The high thermal conductivity of electrons in the direction of the magnetic field then ensures that their temperature is equal to the electron temperature of the plasma outside the trap, but is small in comparison with the value given by (18). Accordingly, since the ion slowing down time τ_{ie}^{E} is proportional to $T_e^{3/2}$, it will be much shorter than the value of τ estimated from (14). Under these conditions, an injected ion is initially slowed by electrons down to energies significantly lower than W_{inj}, and is then rapidly ($\tau_{ii}\sim W_i^{3/2}$) scattered into the loss cone. In other words, the order of magnitude of the ion lifetime is determined by the ion slowing-down time by electrons, and is approximately equal to τ_{ie}^{E}.

Most of the experiments completed so far, have been performed under these conditions. As an example, consider the experiments on the 2XIIB installation,[23] mentioned in the Introduction. The installation is illustrated schematically in Fig. 6a and the measured ion slowing-down time is shown in Fig. 6b. The shape of the plasma boundary is also indicated approximately in Fig. 6. The transverse dimension of the plasma at the center of the system is approximately 20 cm.

2.3. Equilibrium and magnetohydrodynamic (MHD) stability of plasmas

It is clear from the discussion given in Section 2.1 that the plasma in the mirror trap is anisotropic. In a Cartesian system of coordinates, with the z axis pointing along the magnetic field, the momentum flux density tensor is

$$\begin{pmatrix} p_\perp & 0 & 0\\ 0 & p_\perp & 0\\ 0 & 0 & p_\parallel \end{pmatrix}. \tag{19}$$

The off-diagonal elements $p_{\alpha\beta}$ of this tensor are proportional to the ratio of the square of the Larmor radius of ions to the transverse dimension of the plasma (see Ref. 24), and usually have very little effect on equilibrium. We have set them all equal to zero. The quantities $p_\parallel$ and $p_\perp$ are called the longitudinal and transverse plasma pressures, respectively. When atomic beams are injected across the magnetic field in the equatorial plane of the mirror trap, the condition $p_\perp > p_\parallel$ is satisfied throughout the system.

When the off-diagonal elements are neglected in (19), the momentum flux density tensor in an arbitrary coordinate frame has the form

$$p_{\alpha\beta}=p_\perp\left(\delta_{\alpha\beta}-\frac{B_\alpha B_\beta}{B^2}\right)+p_\parallel\frac{B_\alpha B_\beta}{B^2}, \tag{20}$$

where B_α are the components of the magnetic field vector and B is the modulus of this vector. The expression given by (20) resembles the Maxwellian magnetic stress tensor (see Ref. 25)

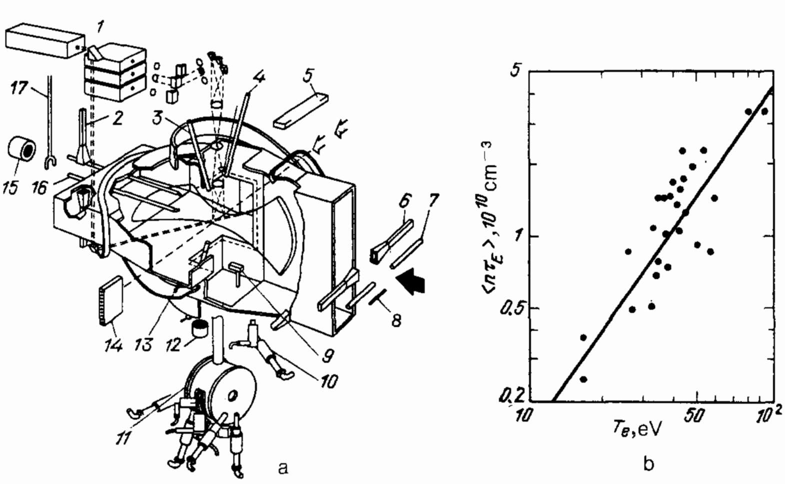

FIG. 6. Experiment on the 2XIIB installation: a—principle of the experiment showing different plasma diagnostic facilities: *1*—Thomson scattering system, *2*—fixed microwave interferometer, *3*—system for detecting microwave scattering at 26°, *4*—microwave interferometer, *5*—neutron counter, *6*—centimeter-band interferometer, *7*—4-mm interferometer, *8*—Langmuir probe, *9*—atomic beam calorimeter, *10*—mobile analyzer of charge-exchange neutrals, *11*—eleven-channel analyzer of charge-exchange neutrals, *12*—x-ray detector, *13*—diamagnetic loop, *14*—thirteen-channel detector of beam attenuation, *15*—electrostatic end-loss analyzer, *16*—mobile 4-mm interferometer, *17*—radio-frequency probe; thick arrow shows the direction of injection of the cold plasma jet; b—energy lifetime of ions $\tau_E\equiv\overline{W}_i|d\overline{W}_i/dt|^{-1}$ as a function of electron temperature; solid line represents the formula $\langle n\tau_E\rangle=4.4\;10^7\;T_T^{3/2}$, where T_e is in electron volts.

$$T_{\alpha\beta}=p_M\left(\delta_{\alpha\beta}-\frac{B_\alpha B_\beta}{B^2}\right)-p_M\frac{B_\alpha B_\beta}{B^2}, \qquad (21)$$

where

$$p_M=\frac{B^2}{8\pi} \qquad (22)$$

is usually referred to in plasma physics as the *magnetic pressure.* Comparison of (20) with (21) shows that the magnetic field exerts a pressure in directions perpendicular to the lines of force, and a tension in the direction of the lines of force. The equilibrium equations for a plasma in a magnetic field can be written in the form

$$\frac{\partial}{\partial x_\alpha}(p_{\alpha\beta}+T_{\alpha\beta})=0. \qquad (23)$$

When applied to the vacuum magnetic field (curl$\mathbf{B}=0$), equation $\partial T_{\alpha\beta}/\partial x_\beta=0$ shows that the force exerted on a particular volume by the magnetic field is zero when the magnetic pressure and stress forces balance. The appearance of the plasma in the magnetic field leads to a distortion of the initial vacuum field and to a force acting on the plasma. Of course, this force is due to the current flowing through the plasma and is equal to

$$\frac{1}{c}[\mathbf{j}\times\mathbf{B}],$$

where **j** is the current density. It is this plasma current that distorts the vacuum magnetic field.

Let us now introduce the dimensionless parameter

$$\beta=\frac{p}{p_M}, \qquad (24)$$

where the magnetic pressure p_M and the plasma pressure p[6] are taken at some characteristic point. It is clear that, in equilibrium, the pressure p cannot be much greater than p_M since otherwise the second term in (23) would be much smaller than the first, and magnetic forces would not balance the gas-kinetic pressure. We may therefore definitely conclude that β cannot exceed unity:

$$\beta\leqslant 1. \qquad (25)$$

The question whether or not β can reach values of the order of unity and, if not, what is the limiting value of β in a given magnetic configuration, requires a quantitative analysis of the equations of equilibrium. This is a very complex problem, especially in non-axially-symmetric cases, and must usually be solved by numerical methods. Moreover, if the problem does not have some particular small parameters associated with the geometry of the system, then it follows from (23) that the limiting value of β will be of the order of unity. Such configurations include "short" mirror traps (Fig. 7a) in which the distance between the mirrors is of the order of the transverse dimension of the plasma. We thus arrive at the conclusion that, from the point of view of plasma equilibrium, there are definitely open-ended traps in which the limiting value of β is of the order of unity. This means that, to confine a plasma at a given pressure p it is sufficient to have a magnetic system capable of producing a magnetic pressure $p_M\sim p$, but not necessarily $p_M\gg p$, as in the case of other plasma confinement systems.

We also note that, according to the equilibrium conditions, the value $\beta\sim 1$ can also be attained in "long" axially symmetric mirror traps in which the transverse dimension of the plasma is small in comparison with the separation between the mirrors (Fig. 7b).

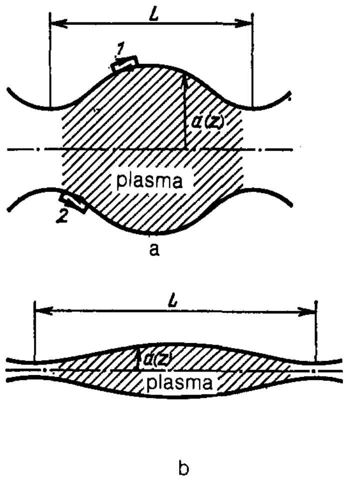

FIG. 7. Short (a) and long (b) magnetic mirror traps: contours 1 and 2 are used to determine whether the magnetic field increases or decreases with distance from the convex and concave plasma boundary, respectively.

Not all the configurations that are acceptable from the point of view of MHD equilibrium are found to be MHD stable. The question of stability is usually examined by considering the change in the potential energy of the system, i.e., the sum of magnetic energy and the internal energy of the plasma for a small deviation of the system from the state of equilibrium. This deviation is described by the function $\xi(\mathbf{r})$, i.e., the displacement of points in the plasma relative to their initial positions. Because of the high plasma conductivity, this motion is accompanied by a perturbation of the magnetic field, which can also be expressed in terms of the function $\xi(\mathbf{r})$:

$$\delta\mathbf{B}=\text{curl}[\xi\times\mathbf{B}]. \qquad (26)$$

As a result, the change $\delta\Pi$ in the potential energy can be written in the form of a quadratic functional of the displacement $\xi(\mathbf{r})$. To prove stability, we have to show that the functional $\delta\Pi$ is positive definite. The proof of instability consists of finding a trial function $\xi(\mathbf{r})$ that makes $\delta\Pi$ negative. Methods of investigating these problems have now reached a high degree of refinement and perfection (see, for example, Ref. 26).

On the whole, the conclusion that can be drawn from such studies is that the plasma tends to flow into the region of weaker magnetic field. This can be understood qualitatively by recalling that the quantity μ is conserved for plasma particles in MHD perturbations[7] and, accordingly, when a particular plasma element is displaced into a region of weaker field, its internal energy is reduced. This statement does not have direct proving power, but it does usually enable us to identify the most hazardous types of perturbation.

Let us now consider the boundary of the plasma confinement region (see Fig. 7a). Evaluating the line integral of the magnetic field over a contour close to the plasma boundary, and noting that this integral is zero in vacuum, we immediately find that the magnetic field decreases with distance from the plasma boundary when this boundary is convex (contour 1) and increases when it is concave (contour 2). In this sense, the plasma boundary curvature can be referred to as favorable or unfavorable.

Plasma deformations that are not accompanied by the distortion of the magnetic field are of particular importance in the problem of MHD plasma stability. It is clear from (26) that deformations for which

$$[\xi \times \mathbf{B}] = \nabla \psi, \tag{27}$$

i.e.,

$$\xi_\perp = \frac{[\mathbf{B} \times \nabla \psi]}{B^2}, \tag{28}$$

where ψ is a certain scalar function and the subscript $\perp$ is used to indicate the direction relative to the magnetic field [longitudinal displacement is not determined by (27)]. For deformations of the form specified by (28), the potential-energy perturbation $\delta\Pi$ consists of only the perturbation of the internal energy of the plasma. Accordingly, the sign of $\delta\Pi$ does not depend on the magnetic field. In particular, for an inappropriate choice of the magnetic-field configuration, $\delta\Pi$ can be negative even for $\beta \to 0$, i.e., even very low pressure plasma will be unstable. This ensures that perturbations such as (28) occupy a special place. Their instability means that plasma with appreciable pressure cannot accumulate in this system. Such perturbations are called *interchange perturbations* because they can be looked upon as the result of the interchange of tubes of force together with the plasma contained by them (Fig. 8). They are also referred to as *flute perturbations.*

The stability of plasma with respect to interchange perturbations was investigated as far back as the 1950s, in connection with open-ended traps.[27,28] The stability condition for axially-symmetric plasma was obtained by Rosenbluth and Longmire.[27] It takes the following very simple form in the so-called paraxial limit, i.e., in plasma whose radius a is small in comparison with the characteristic scale L of changes in the magnetic field along the axis of the system:

$$\int \frac{p_\parallel + p_\perp}{B_0(z)^{3/2}} \frac{d^2 a}{dz^2} dz > 0, \tag{29}$$

where $B_0(z)$ is the magnetic field on the axis of the system and $a(z)$ is the distance between the axis and the plasma boundary (coincident with one of the lines of force). It is assumed that the plasma inside the boundary is uniform over the cross section. The integral is evaluated between the mirrors (where the plasma pressure is zero).

In the paraxial approximation, the quantity da^2/dz^2 is equal to the curvature of the plasma boundary. Condition (29) is a quantitative expression of the previously-mentioned role of favorable and unfavorable curvatures.

If we denote the radius of the plasma and the magnetic field in the equatorial plane of the trap by a^* and B_0^*, respectively, then the equation of the plasma boundary in the paraxial approximation can be written in the form

$$a^2 B_0 = a^{*2} B_0^*.$$

Let us now express a in terms of $\mathbf{B}_0$, and substitute the result in (29). The plasma pressure on a line of force can be looked upon as a function of the magnetic field: $p_\parallel = p_\parallel(B_0)$, $p_\perp = p_\perp(B_0)$. Bearing all this in mind, and evaluating the integral in (29) by parts, we obtain the following inequality instead of (29):

$$\int \frac{dz}{B_0^{3/2}} \left(\frac{dB_0}{dz}\right)^2 \frac{d}{dB_0} \frac{p_\parallel + p_\perp}{B_0^{3/2}} > 0. \tag{30}$$

Under "standard" conditions, in which particles are injected into the trap in the equatorial plane across the magnetic field, the quantities $p_\parallel$ and $p_\perp$ are decreasing functions of B_0, i.e., the plasma pressure falls in the direction of the mirrors. Condition (30) then shows that the plasma in the axially-symmetric mirror trap is unstable with respect to interchange perturbations.

To ensure the MHD plasma stability, it is desirable to have a magnetic field configuration in which all the lines of force are always convex to the plasma. The required effect can be achieved by abandoning the axial symmetry of the problem. Figure 9 shows a magnetic field configuration that has the necessary property. It is obtained from an initially cylindrical tube of lines of force by flattening the ends of the tube in mutually perpendicular directions. In these configurations, the lines of force are convex to the plasma, and the magnetic field always increases towards the periphery of the plasma. They are referred to as "minimum-B" configurations.

The possibility of plasma stabilization by minimum-B configurations was first demonstrated experimentally by the group headed by M. S. Ioffe[3] in the early 1960s.[8)]

When the mirror ratio R is large, ($R \gg 1$), the plasma can be almost isotropic in most of the trap. This type of plasma can be characterized by a scalar pressure $p(p_\parallel \approx p_\perp)$. Projecting the equations of equilibrium (23) on to a line of

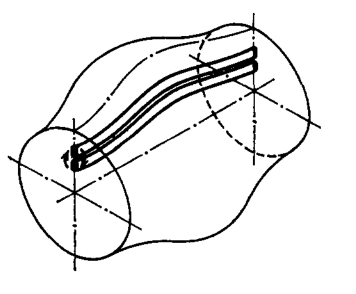

FIG. 8. Flute (interchange) perturbations.

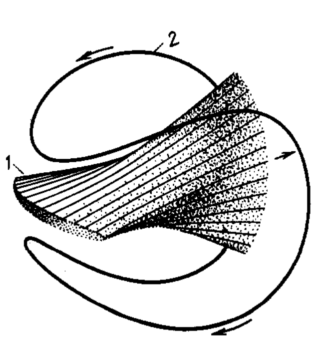

FIG. 9. "Minimum-B" magnetic field configuration (from: *The National Mirror Fussion Plan,* Livermore, 1980). The figure shows one of the magnetic surfaces with its individual lines of force (1). The coil producing the required field configuration is also indicated (2). Because of the quadrupole symmetry properties of the magnetic potential of this configuration it is often referred to as a *quadrupole.*

force, we readily see that the pressure is then constant along the line of force (of course, with exception of the narrow region near the mirrors, in which the plasma is anisotropic and the pressure falls to zero). The pressure is thus a function of the line of force.

A stability criterion that is free from the assumption that the system is axially symmetric was derived by B. B. Kadomtsev[28] for isotropic plasma. The quantity

$$U = \int \frac{ds}{B}, \tag{31}$$

plays an important part in this criterion, where the integral is evaluated along a line of force between the mirrors (we recall that these mirrors are assumed to be very strong). The quantity U is also a function of the line of force. Its significance can be established by multiplying it by $\delta\Phi$, i.e., the magnetic flux through a narrow tube of force containing the chosen line of force. Since $\delta\Phi/B$ is the cross sectional area of a tube, it follows that $U\delta\Phi$ is the volume δV of the tube enclosing the flux $\delta\Phi$, i.e., U can be interpreted as the *specific volume* of the tube ($U = \delta V/\delta\Phi$).

It can be shown that, in equilibrium, the plasma pressure is constant over the surface $U =$ const, i.e., $p = p(U)$.[28] We shall label each magnetic surface $U =$ const with the value of the flux Φ enclosed by it. In this sense, we can speak of the function $U = U(\Phi)$ and, similarly, of the function $p = p(\Phi)$.

In terms of the functions $U(\Phi)$ and $p(\Phi)$, the stability criterion obtained by Kadomtsev takes the form

$$-\frac{dp}{d\Phi}\frac{dU}{d\Phi} < \frac{5p}{3U}\left(\frac{dU}{d\Phi}\right)^2. \tag{32}$$

Since for the magnetic field shown in Fig. 9, the field B increases sufficiently rapidly with distance from the magnetic axis, we have $U'_\Phi < 0$ (at least at distances from the magnetic axis that are not too large). Accordingly, for the plasma pressure profiles with $p'_\Phi < 0$, which are of interest in relation to magnetic plasma confinement, the stability condition (32) is definitely satisfied.

The function $U(\Phi)$ changes by an amount of the order unity when the distance from the axis increases from zero by an amount of the order of the separation L between the mirrors. Hence, if the plasma fills only the paraxial region with a transverse size $a \ll L$, then $|p'_\Phi|/p \gg |U'_\Phi|/U$. Accordingly, the right-hand side of inequality (32) can be replaced with zero in the paraxial limit, and the stability condition can be written in the form $p'_\Phi U'_\Phi > 0$ which, for $p'_\Phi < 0$, reduces simply to

$$U'_\Phi < 0. \tag{33}$$

Of course, in the axially-symmetric case, conditions (33) and (29) become equivalent to one another (for $p_\parallel + p_\perp = 2p =$ const) (the proof requires some further algebra).

When the mirror trap is intrinsically unstable, it can be stablized by providing it at each end with an additional mirror trap that has a sufficiently deep "minimum-B" [so that condition (33) is satisfied for the system as a whole]. The system is then said to exhibit a "mean minimum-B" (mean along a line of force), and the stabilizing traps are referred to as "anchors" or "*stablizers.*"

If the system is stable with respect to flute perturbations, it must be checked for stability with respect to all other trial functions $\xi(\mathbf{r})$. When the values of β are high enough, perturbations similar to flute perturbations, but localized in the direction of the magnetic field in the region of unfavorable curvature of the lines of force, i.e., the so-called "balloon" perturbations, may be energetically favorable. The resulting limits on β can be more stringent than those derived from the equilibrium conditions. However, detailed analysis of this problem has shown that the limiting value of β decreases only slightly when the magnetic configuration is suitably chosen.

"Long" mirror traps (in the sense of Fig. 7b) exhibit an important stabilizing effect, predicted by Rosenbluth, Krall, and Rostoker[29] and due to the finite Larmor radius (FLR) of ions. The effect is caused by the off-diagonal terms of the tensor $p_{\alpha\beta}$ that were neglected in (19). They are proportional to r_{Li}^2 and are similar in form (see Ref. 24) to the elements of the viscous stress tensor (although they do not lead to the dissipation of the perturbation energy). When the plasma radius a is sufficiently small in comparison with the length L, this *nondissipative viscosity* prevents the development of interchange perturbations, except for those with the largest scales, which correspond to the displacement of the plasma as a whole in the direction perpendicular to the magnetic axis. The condition under which small-scale perturbations are stablized by the FLR effect is

$$Lr_{Li} \gg a^2.$$

From the standpoint of MHD stabilization, there is some interest in the magnetic-field configuration due to two coaxial mirrors connected in opposition and producing field lines as shown in Fig. 10. This is the so-called cusp (*antiprobkotron* in Russian literature) in which the lines of force are convex to the confinement region and ensure the MHD plasma stability. Unfortunately, the confinement region contains the zero-field point O. The adiabatic invariant μ is then no longer conserved along the lines of force passing near O, and there is a rapid loss of particles along the lines of force. This results in an empty space inside the confinement

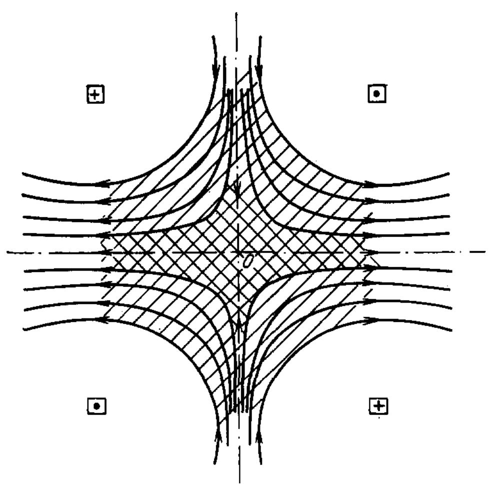

FIG. 10. The cusp: arrows show the magnetic lines of force and O is the point at which the magnetic field is zero; the region occupied by the plasma is cross hatched; double cross hatching shows the region from which particles are lost along the lines of force due to the violation of the adiabatic invariance of μ when particles travel near the zero field point.

region (see Fig. 10), and the boundary of the ambient plasma is convex to this space, i.e., it is unstable. The instability can lead to the filling of the neighborhood of the null point, the escape of plasma along the lines of force, the onset of new instability, and so on, until there is no more plasma in the trap.

The longitudinal loss of plasma out of the nonadiabatic region must be prevented if the cusp configuration is to be used for the purposes of controlled thermonuclear fusion. This point is discussed in Section 4.

2.4. Microinstabilities

The"hole" in velocity space that is not filled with ions can serve as a source of free energy for the spontaneous excitation of different types of short-wave plasma oscillations (with characteristic wavelength equal to, or less than, the Larmor ion radius). The interaction between the plasma ions and the electromagnetic field of such oscillations will, in general, produce a random variation in the magnetic moment μ and a rapid scattering of ions into the loss cone.

The number of potentially hazardous types of oscillation is very large, as is the number of possible mechanisms that can excite them. Traditionally, these mechanism are commonly referred to as microinstabilities.

One of the most hazardous is the so-called drift-cone instability (DCI) that was predicted and theoretically investigated in Refs.30 and 31. This is excited in the mirror-trap plasma when the radius of the plasma column is small enough (more precisely, the radial density gradient is high enough). The perturbations take the form of "flutes" elongated along the magnetic lines of force and running azimuthally in the direction of the diamagnetic drift of ions. The characteristic frequency of the perturbations is of the order of ion cyclotron frequency ω_{Bi}.

Since DCI can develop even in low pressure plasmas ($\beta \gg 1$), and it originates in the universal properties of the mirror-trap plasma such as the loss cone (or, more precisely, the loss hyperboloid) and the radial density inhomogeneity, this instability could have finally sealed the fate of the magnetic mirror trap.

Fortunately, as was noted in Ref.31, the instability vanishes when a small number of ions appears near the $v = 0$ point. These ions are referred to as "warm," in contrast to the "hot" ions of the main plasma.

Of course, the"warm" ions are not confined within the trap [see (12) and Fig. 5], and certain additional measures must be taken to maintain their population. In the experiments of Ioffe *et al.*, this was done by producing a shallow "well" in the electrostatic potential in the central part of the trap, which ensured the electrostatic confinement of slow ions (see Sec. 3.5 for further details). A continuous-flow arrangement was used in the 2XIIB system, whereby warm plasma was introduced into the system through one tube and was removed through another. Since the plasma outside the trap was in direct contact with the end walls of the system, its presence in the system caused a cooling of electrons, i.e., the electron temperature was restricted to 50–100 eV, whereas the energy of the injected protons was 20 keV. The electron temperature could not be raised by reducing the rate of plasma inflow because DCI developed in the system and the loss of fast ions became unacceptably high.

As the plasma radius increases (radial pressure gradient decreases), the amount of warm ions necessary for stabilization decreases and becomes quite small as we pass to parameter values typical for reactor plasma,[32] so that the problem of DCI stabilization in the mirror-trap reactor can probably be solved (although the problem of maintaining a stationary population of warm ions in the plasma of a mirror-trap reactor is not at all simple even when this population is quite low).

As already noted, the number of potentially hazardous microinstabilities in the mirror trap is very large. The difficulty of providing a satisfactory theoretical description of these instabilities is that the problem has a large number of parameters, and the plasma and magnetic field are often longitudinally and transversely nonuniform. Nevertheless, the linear theory (i.e., in problems of the evolution of small initial perturbations) has attained a considerable degree of clarity, but not in all areas by far (see review papers of Refs. 19, 26 and 32). However, in the final analysis, the role of particular instabilities is determined by the level reached by the corresponding fluctuations and by the effective time τ_{eff} for the scattering of ions by these fluctuations. The answers to such questions can only be provided by the *nonlinear* theory which is, clearly, still incomplete. It follows that experiments satisfying the necessary similarity conditions (in relation to the thermonuclear reactor) must provide the final answer. The basic similarity parameters are

$$N \equiv \frac{a}{r_{Li}},\ \frac{a}{L},\ \frac{\omega_{Bi}}{\omega_{pi}},\ \beta,\ \frac{T_e}{W_{inj}}, \tag{34}$$

where a is the plasma radius, L is its length, and ω_{pi} the plasma ion frequency. The remaining notation is the same as above. In addition, the ion Coulomb scattering time τ_{ii} [see (7′)] should be made significantly greater than the time for charge exchange with the residual gas.

Unfortunately, experiments performed so far do not satisfy these conditions, and no final conclusion can be made about the significance of the conditions for "classical" (i.e., unrelated to scattering by microfluctuations) plasma confinement in the mirror-trap reactor. We merely note that the requirement of "classical" confinement imposes stringent restrictions on the admissible level of microfluctuations. As an illustration, consider the scattering of ions by microfluctuations with characteristic scale of the order of the Debye radius r_D. We know (see, for example, Ref. 33) that the effective time for scattering of ions by microfluctuations, τ_{eff}, can be estimated from the formula

$$\tau_{eff} \sim \frac{\ln N_D}{N_D}\,\frac{nT_i}{\mathscr{W}},$$

where $\mathscr{W}$ is the microfluctuation energy density and N_D is the number of particles in the Debye sphere. The condition $\tau_{eff} > \tau_{ii}$ gives the following limitation on the admissible level of fluctuations. This is an unusually stringent restriction.

$$\frac{\mathscr{W}}{nT} < \frac{\ln N_D}{N_D}.$$

In fact, for a thermonuclear fusion plasma with $n \sim 10^{14}$ cm^{-3}, $T_i \sim 2\cdot 10^5$ eV, we have $N_D \sim 5\cdot 10^{10}$, i.e., the ratio $\mathscr{W}/nT$ must be less than 10^{-9}. This shows the extent to which plasma must be stable with respect to these microfluctuations.

To gain some idea about the possible influence of microfluctuations on the energy parameters of an open-ended trap, it is useful to investigate the dependence of these parameters on the quantity defined by

$$\mathscr{P} \equiv 1 + \frac{\tau_{ii}}{\tau_{eff}} > 1, \tag{35}$$

where τ_{ii} is the Coulomb scattering time (7'). The part played by $\mathscr{P}$ in this analysis is such that it may be referred to as the "gloom factor." Since it is probable that epithermal microfluctuations in the open-trap plasma cannot be suppressed, it is unlikely that systems requiring a gloom factor $\mathscr{P} < 1.5$–2 for energy balance will function normally as thermonuclear reactors.

2.5. The mirror trap as a thermonuclear reactor

In the near future, the fuel for fusion reactors will be a mixture of deuterium and tritium because the cross section for the reaction

$$D + T \rightarrow He^4 + n \tag{36}$$

is much greater than the cross sections for the other fusion reactions ($D + D \rightarrow He^3 + n$ and $D + He \rightarrow He^4 + p$). The energy yield of (36) is $W_f = 17.6$ MeV, and this is divided in the ratio of 4:1 between the neutron and the α particle.

The fusion power produced per unit volume in a mixture of equal amounts of deuterium and tritium is

$$P_f = \frac{1}{4} n^2 W_f \langle \sigma_{DT} v_{DT} \rangle,$$

where n is the total number of D and T nuclei per unit volume, σ_{DT} is the cross section for the reaction (36), v_{DT} is the relative velocity of the D and T nuclei, and the angle brackets represent averaging over the deuteron and triton distribution functions.

The external injection sources supply the plasma with the energy and matter necessary to compensate for losses through the mirrors. In the steady state, the power injected per unit plasma volume should be[9]

$$P_{inj} = \frac{n W_{inj}}{\tau},$$

where τ is the ion lifetime in the trap. Strictly speaking, the heating of electrons by α particles does provide some contribution to the energy balance but, as we shall see, this is relatively insignificant in mirror traps.

The energy multiplication factor of a fusion reactor is defined by

$$Q = \frac{P_f}{P_{inj}} = \frac{1}{4} \frac{W_f}{W_{inj}} n\tau\sigma \langle_{DT} v_{DT} \rangle.$$

For a given magnetic-field geometry (in particular, a given mirror ratio), the product $n\tau$ and the quantity $\langle \sigma_{DT} v_{DT} \rangle$ are functions of only the injection energy, and do not depend on plasma density. Numerical calculations that are reviewed, for example, in Refs. 16 and 19, show that Q increases with injection energy up to $W_{inj} = 200$–300 keV, and then very slowly decreases. The maximum value of Q even for $R = 10$ is only 1.4, and the logarithmic dependence of τ on R [see (14)] prevents us from producing a significant increase in Q by increasing R[10]. Thus, even under the boldest assumptions about the possible values of R, one can hardly expect to achieve energy multiplication factors exceeding 1.5–2 in mirror traps. It is precisely this conclusion that is the basis for the rather gloomy view of the mirror trap as a possible reactor. The point is that, when a realistic view is taken of the efficiency of conversion of thermonuclear energy (released mostly in the form of the neutron flux) into the energy of deuterium and tritium beams injected in the plasma, a "closed" energy balance cannot be established in the system for $Q < 3$–4.

Moreover, we have also made the optimistic assumption that microfluctuations will not lead to a reduction in the lifetime as compared with the estimate given by (14), which is likely not justified (see Sec. 2.4).

Let us summarize the situation so far. Many of the properties of mirror traps suggest that it is a very successful fusion system: it makes effective use of the confining magnetic field ($\beta \sim 1$ can be reached), it can operate under steady conditions (there is no problem with the build-up of reaction products), it is not very sensitive to the presence of impurities, and it is "topologically" simpler than closed systems. However, all these advantages are paid for by the fact that high energy multiplication factors cannot be achieved in such traps.

In the following Sections, we shall consider the extent to which this defect can be removed while retaining the advantages of the mirror trap.

3. AMBIPOLAR TRAP

3.1. Principles

The advent of the concept of the ambipolar trap was one of the most sensational events in the history of research on controlled thermonuclear fusion: very simple (at least in principle) and well established (individually) procedures were suggested as a means of improving the simple mirror trap, and this has reinstated the question of the open-ended mirror trap as a reactor. Moreover, as stated in the Introduction, the idea of the ambipolar trap fell on well-prepared ground: the experimental techniques necessary for its implementation had already reached the necessary level of development, and had been tested on the 2XIIB at Livermore.

The ambipolar trap was first described by D. Y. Dimov, V. V. Zakaĭdakov, and M. E. Kishenevskiĭ in a paper published[11] in 1976. The paper by T. K. Fowler and E. G. Logan reporting a similar proposal was published[12] at the beginning of 1977. The principle of the ambipolar trap can be explained as follows. Consider a mirror trap (2 in Fig. 11) with an additional mirror trap added at each end (1 and 3 in Fig. 11). Suppose that high-density plasma is maintained in these end traps by the intensive injection of atomic beams (this, of course, requires considerable expenditure of energy). Let us consider how the presence of dense plasma in the end cells 1 and 3 affects the confinement of ions in the central cell 2. We recall that, because of the high frequency of electron–electron collisions (see Sec. 2.1), the electron distribution function is Maxwellian and, that to ensure that the local electron and ion densities are equal, the ambipolar potential whose distribution along the line of force is given by (9) is established in the plasma. In the present case, the potential distribution has the form shown in Fig. 11, i.e., ions in the central cell are located in a potential well of depth

$$e\Delta\varphi = T_e \ln \frac{n_{max}}{n_0}, \tag{37}$$

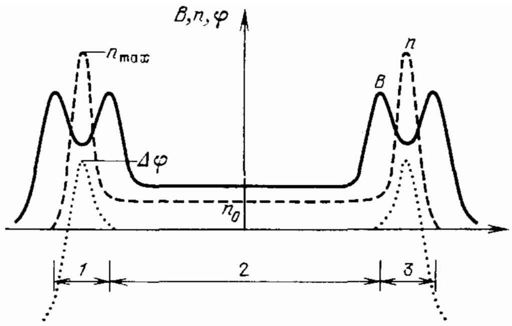

FIG. 11. Principle of the ambipolar trap: the distributions of magnetic field (solid line), plasma density (dashed line), and electrostatic potential (dotted line) are shown along the length of the system; numbers label the central (2) and the two end cells (1, 3).

where $\Delta\varphi$, n_{max}, and n_0 are defined in Fig. 11. Ions of energy less than $e\Delta\varphi$ are confined to the central cell whatever the direction of their velocity vector. When $e\Delta\varphi \gtrsim T_i$, the ion distribution function is nearly Maxwellian, and only ions in the Maxwellian tail can leave the trap. Accordingly, the ion lifetime rises approximately by the factor

$$\exp\frac{e\Delta\varphi}{T_i} = \left(\frac{n_{max}}{n_0}\right)^{T_e/T_i}$$

as compared with (14).[11] In principle, the ion lifetime in the central cell can be made as large as desired. Of course, as already noted, the high plasma density in the end traps is achieved at the cost of some expenditure of energy but, by increasing the length of the central cell, it is always possible to ensure that the fusion energy released in the central cell is much greater than the energy input into the end cells (which, we emphasize, is independent of the length of the central trap).

This is the general idea of the ambipolar trap. It is interesting to note that a system of three mirror traps was considered by G. G. Kelley[35] as far back as the middle 1960s. Kelley proposed the use of small additional mirror traps, with roughly the same density as in the main trap, as a means of reducing the undesirable effect of the extracting ambipolar potential on ion confinement in the main trap. Almost ten years elapsed before it was noted that the relatively passive end traps could be given a controlling and actively positive role, whereby a radical improvement could be made in the energy parameters of the system. It is here, and not in the replacement of one trap with a set of three traps (as it is sometimes stated), that we see the essence of the ambipolar trap.

Although the energy balance in the ambipolar trap can be established, at least in principle, for an arbitrarily high rate of injection of energy into the end traps, it is desirable to keep this injected power as low as possible in order to ensure that the central cell is not too long (say, not more than 1 km). For a given plasma volume and density in the end cell, the loss of power from it (and, correspondingly, the injected power) is proportional to $1/W_{inj}^{1/2}$ [see (14) and (7')]. Accordingly, the injected energy must be chosen to be as high as possible. In the first publication[11] on ambipolar traps, it was assumed that W_{inj} lay in the range 1–2 MeV. This meant that a very strong magnetic field had to be employed in the end traps, since otherwise their size would have been unacceptably large. A field of 200 kG was considered (for the mirrors) in Ref. 11. The length of the central cell was found to lie in the range between a few hundred meters and 1 km.

When ion losses from the end traps are determined not by Coulomb collisions but by the scattering of ions by microfluctuations [this is indicated by the gloom factor $\mathscr{P} \gtrsim 1$ in (35)], the amount of energy that has to be expended in maintaining the high plasma density in end mirrors is higher. However, in contrast to the simple mirror trap, this energy expenditure can now be compensated by increasing the length of the central cell. Of course, if the coefficient $\mathscr{P}$ becomes too high, say, $\mathscr{P} > 20$–30, the required length of the central cell becomes unreasonable. Technical restrictions on the attainable specific injected power (power per unit plasma volume) are still more stringent.

3.2. Experimental verification of ambipolar confinement

The ambipolar confinement effect was demonstrated experimentally on Gamma 6 at Tsukuba[36] and the TMX at Livermore.[37] We shall confine our attention to the results obtained on the second of these two machines.

The magnetic field structure of the ambipolar trap used in the TMX is illustrated in Fig. 12, which shows one of the magnetic surfaces. The stability of the system with respect to flute perturbations is assured by the end mirror traps in which there is a deep "minimum-B." Thus, the end mirror traps not only ensure that the plasma is confined to the center cell, but they also serve as MHD stabilizers (see Sec. 2.3).

The end mirrors are turned through 90° about magnetic axis of the system relative to one another. The point of this is that the $U = \text{const}$[12)] surfaces (which coincide with constant plasma-pressure surfaces) are very nearly circular cylinders over a long portion of the uniform magnetic field (as shown in Fig. 12). For a different relative orientation of the end mirrors, the $U = \text{const}$ surfaces would not in general be circular cylinders, which would be undesirable for a number of reasons (in particular, it would produce a greater ratio of surface area to plasma volume).

The basic parameters of the system are as follows: distance between center-cell mirrors 6.4 m, length of each of the end cells 0.75 m, magnetic field in the mirrors 20 kG, beam injection power 3.5 MW (in each of the end cells), and injection energy 20 keV. The machine as a whole is a complex and

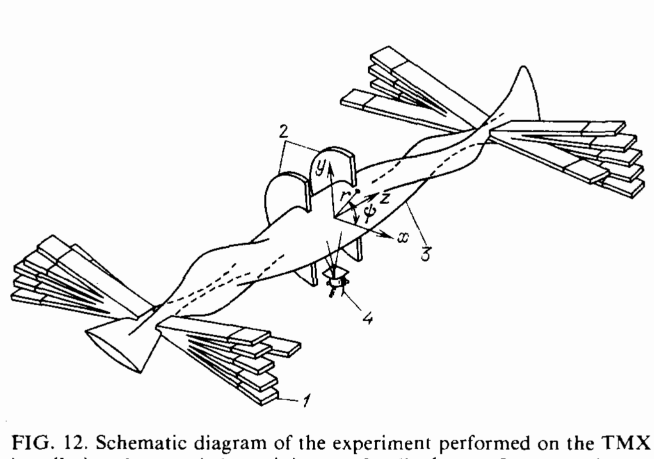

FIG. 12. Schematic diagram of the experiment performed on the TMX installation: 1—atomic beam injectors, 2—diaphrams, 3—magnetic surface bounding the plasma, 4—gas admission valve.

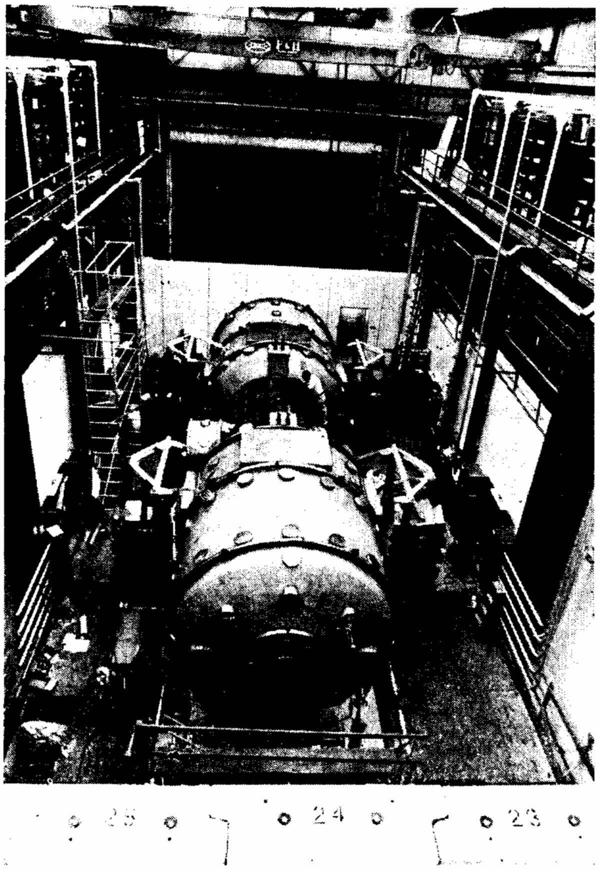

FIG. 13. Photograph of the TMX installation (kindly supplied by F. H. Coensgen).

highly sophisticated engineering construction. A photograph of it is shown in Fig. 13.

The plasma parameters depend on the power injected into the end plugs, the amount of gas supplied to the central cell, and other experimental conditions. Typical parameter values were as follows: plasma density in the central cell 5×10^{12} cm^{-3}, temperature of "central" ions 60 eV, plasma density in end cells 2×10^{13} cm^{-3}, mean ion energy 13 keV, and electron temperature 100 eV.

By varying the injection current into the end cells and the supply of gas for the central cell, the ratio n_{max}/n_0 can be varied between wide limits (see Fig. 11 for the notation). It has been found that an increase in this ratio by a factor of about three led to an increase in the lifetime in the central cell by a substantial factor (up to nine) as compared with the case where there was no ambipolar confinement. The lifetime was in satisfactory agreement with the Pastukhov formula mentioned in Sec. 3.1. The feasibility of ambipolar traps has therefore been fully confirmed experimentally.

On the other hand, a further increase in the ratio n_{max}/n_0 gives rise to plasma microfluctuations with frequencies approaching ω_{Bi}, and the lifetime begins to fall. The fluctuations are nonpotential (they are accompanied by appreciable magnetic field perturbations) and are different from those excited in DCI (see Sec. 2.4). They are closer to the Alfven fluctuations due to the anisotropy of plasma pressure and the finite value of β. The complete picture is relatively complicated, and a complete theory is still lacking.

3.3. Loss of plasma across the magnetic field

Soon after the first publications on ambipolar traps, it was shown theoretically[13,38] that a combination of the two factors typical for the central cell of the ambipolar trap, namely, its large length and non-axially-symmetric end plugs, can lead to considerable plasma loss across the magnetic field.

To exhibit the mechanism responsible for this loss, let us consider the drift of an individual ion in the central cell. The ion drifts around the magnetic axis as it travels along the magnetic field produced by the central solenoid. There are two sources of this drift, namely, the radial inhomogeneity of the confining field, due to the finite plasma pressure, and the radial electric field, which is always present in the open-trap plasma. The drift velocity is (see Ref. 39)

$$v_d = \frac{c\mu}{eB}\frac{\partial B}{\partial r} + \frac{c}{B}\frac{\partial \varphi}{\partial r}. \tag{38}$$

Because of this drift, an ion traveling between mirrors is displaced azimuthally through the angle (Fig. 11)

$$\Delta\psi(v_\parallel, v_\perp, r) = \frac{v_d L}{v_\parallel r}, \tag{39}$$

where L is solenoid length, r is the radius of the magnetic surface along which the ion drifts, and $v_\parallel$ is the velocity of the ion along the magnetic field. Numerical estimates show that, under the conditions typical for the fusion reactor, the large values of L ensure that $\Delta\psi$ lies somewhere in the range between 1 and 10 (in the conventional mirror trap, $\Delta\psi \ll 1$).

We now turn to the motion of ions inside the non-axially-symmetric plug. Inside the plug, the curvature vector $\varkappa$ of a line of force will in general have a nonzero tangential component $\varkappa_t$ along the magnetic surface (in the axially-symmetric trap, the lines of force lie in the meridional planes of the trap, and $\varkappa_t = 0$). Accordingly, as the ion travels within the plug, its drift velocity has the component (see Ref. 39)

$$v_n = \mathbf{n}\cdot[\mathbf{B}\times\boldsymbol{\varkappa}]\frac{m_i c[v_\parallel^2 + (v_\perp^2/2)]}{eB^2}, \tag{40}$$

normal to the magnetic surface. As a result, when it is reflected from the asymmetric plug, the ion is displaced by the amount δr relative to the original magnetic surface.[13] This displacement depends on the polar coordinates (r,ψ) of the point of entry of the ion into the plug, and on the values of $v_\parallel, v_\perp$ within the uniform region:

$$\delta r = \delta r(r, \psi, v_\parallel, v_\perp).$$

It can be shown (see Ref. 40) that, in the paraxial approximation,

$$\delta r = \eta r \sin 2\psi + \zeta r^3 \sin 4\psi, \tag{41}$$

where

$$\eta \sim \frac{r_{L1}}{l_p}, \tag{41'}$$

$$\zeta \sim \frac{r_{L1}}{l_p^3} \tag{41''}$$

and l_p is the length of the plug. Since, in practice, we usually have $r/l_p \lesssim 0.1$, only the first term need usually be retained in (41). The estimate (41') is readily obtained from (40) by recalling that, in the paraxial approximation $\varkappa \sim r/l_p^2$, and during the motion of the ion in the plug $\sim l/v_p$.

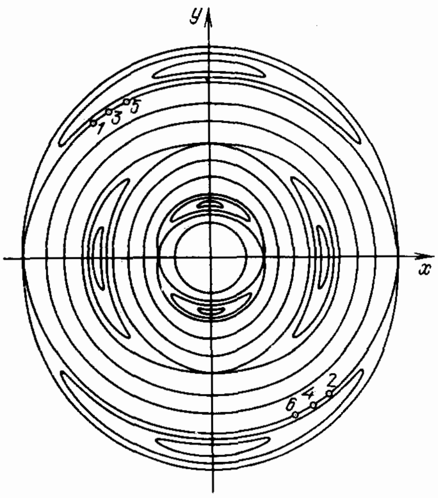

FIG. 14. Family of "resonance" particle trajectories with given energy and magnetic moment. The figure shows successive points at which the drift particle trajectories cross the equatorial plane at time intervals equal to $2t_{\parallel}$.

Since $\langle \delta r \rangle_{\psi} = 0$, the radial displacement averaged over a few periods of longitudinal oscillation is small for most of the particles. There is, however, a class of particle for which the displacements are added constructively on successive reflections from the plugs. These are the so-called resonance particles,[13] i.e., particles which turn through an angle $\Delta\psi$ during their motion between mirrors, so that their reflection occurs at identical points on opposite plugs. For traps with quadrupole plugs turned through 90° relative to one another (see Fig. 12), this condition is

$$\Delta\psi = (2k+1)\frac{\pi}{2} + \varepsilon \quad (k = 0, \pm 1, \ldots), \tag{42}$$

where ε is a small "detuning." If in every fourth transit of a resonance particle across the equatorial plane of the trap we record the point at which the particle crosses this plane, we obtain the quasiperiodic trajectory shown in Fig. 14. The characteristic radial "swing" of the trajectory is given by the following order-of-magnitude formula:

$$\Delta r \sim \left(\frac{\delta r}{|\partial\Delta\psi/\partial r|}\right)^{1/2} \sim \left(\frac{r\delta r}{|\Delta\psi|}\right)^{1/2}$$

and is much greater than δr.

The shape of the resonance trajectories crossing a given point in space depends on the values of $v_{\parallel}$ and $v_{\perp}$. It follows that the Coulomb scattering of resonance particles leads to their "jumping" from one trajectory to another, and this is accompanied by random displacements along the radius by amounts of the order of Δr. The resulting ion diffusion mechanism is often referred to as the resonance mechanism. It is clear that resonance diffusion is actually due to the superposition of two factors, namely, the large length of the central cell [which means that $\Delta\psi$ can reach values exceeding unity and the resonance condition (42) can be satisfied] and the non-axially-symmetric nature of the mirrors (which ensures that δr is finite).

Since the resonance regions are small in velocity space [because the detuning ε in (42) must be small], the frequency of these "jumps" from one resonance trajectory to another is much greater than the quantity τ_{ii}^{-1} [see (7')] that

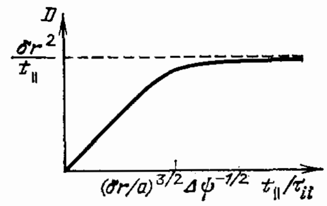

FIG. 15. Diffusion coefficient as a function of collision frequency under resonance conditions.

determines the frequency of ion scattering through an angle of the order of unity. The situation is analogous to that in the "neoclassical" theory of transport processes in tokamaks (see the review in Ref. 11). Accordingly, the dependence of the diffusion coefficient on the ion-ion collisions frequency is also very similar (Fig. 15): for low values of τ_{ii}^{-1}, we have the analog of the "banana" diffusion regime, whereas for large values we have the analog of the plateau regime (see Ref. 41).

The measured rate of transverse loss in the ambipolar traps of TMX and TMX-U at Livermore (USA) and Gamma 10 at Tsukuba (Japan)[42,43] is in satisfactory agreement with the theory of resonance transport. As an illustration, Fig. 16 shows the ion lifetime in the central cell of TMX-U (for escape across the magnetic field) as a function of the potential difference $\Delta\varphi$ between the axial region and the chamber walls ($\Delta\varphi$ is a measure of the radial electric field in the plasma and, correspondingly, of the twist angle $\Delta\psi$). The solid line represents calculations based on the theory of resonance transport.

Future thermonuclear fusion reactors should also work under the resonance transport conditions. Numerical calculations show that, unless special measures are taken to minimize δr, the radial loss may be very considerable. The question therefore arises as to whether δr should be reduced relative to the very approximate estimate given by (41'). The most radical solution is to find fields in which the condition $\mathbf{n}\cdot\mathbf{B}\times\boldsymbol{\varkappa} = 0$ is satisfied throughout the region [see (40)]. Such fields are often referred to as omnigenic. They have the property that particles with arbitrary $v_{\parallel}, v_{\perp}$ that are emitted from the same point will drift over the same magnetic surface. It is clear that the above transport mechanism does not apply to them. A trivial example of an omnigenic field is the axially symmetric field. Nontrivial examples that have ac-

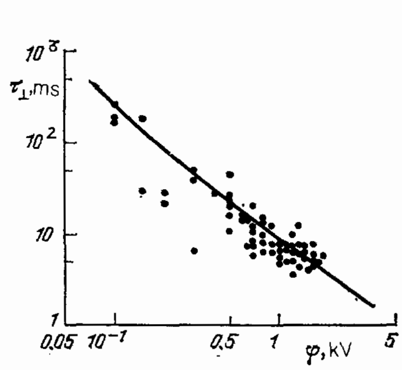

FIG. 16. Ion lifetime in the central cell of the TMX-U installation as a function of the transverse potential difference. The solid line was calculated from the neoclassical resonance formulas.

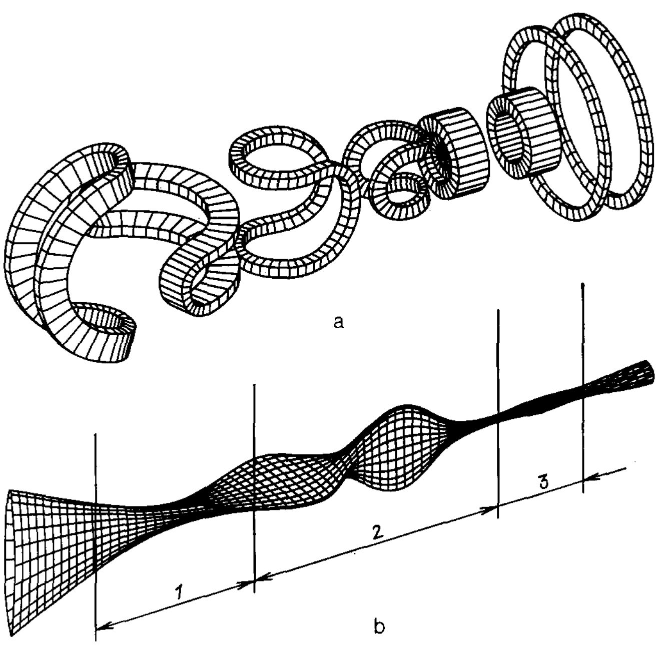

FIG. 17. Magnetic system (a) and magnetic surface (b) of the end section of the MFTF-B installation (Livermore, USA): 1—closing end mirror trap, coincident with the MHD stabilizer, 2—transition region, 3—axially symmetric mirror trap. The central solenoid (not shown) is located on the right.

ceptable shapes in the confinement region and ensure MHD stability have not as yet been constructed. There are arguments suggesting that, most probably, there are no such fields. However, there are fields in which the condition $\mathbf{n}\cdot\mathbf{B}\times\varkappa$ is satisfied approximately: a departure of $\mathbf{n}\cdot\mathbf{B}\times\varkappa$ from zero is then of the third order of small quantities in the paraxial parameter r/l_p (see Refs. 44–46). Unfortunately, attempts to implement even this "approximate" omnigenic property (while retaining MHD stability) with the help of real coils has not led to technologically acceptable solutions either.

Nevertheless, by trying different coil shapes and current distributions in the coils, it is possible to find configurations in which δr is several times smaller than in unoptimized configurations, and thus to achieve reduction of resonance transport coefficients to an acceptable level. An example of a magnetic system found in this way is shown in Fig. 17.

On the whole, the conclusion that can be drawn from dozens of attempts to reduce resonance transport in traps with quadrupole stabilizers is that radial losses can indeed be reduced to an acceptable level, but only by using such complicated magnetic configurations, and at such high cost of having to make them consistent with the conditions for MHD stability and equilibrium infinite-β plasmas, that it would be extremely desirable to find completely axially-symmetric (MHD stable!) magnetic configurations. The solution of this problem would substantially simplify the construction of ambipolar traps, and would enable us to wash our hands of the problem of microinstability suppression.

Finally, let us consider radial electron transport. Since, at temperatures comparable with the ion temperature, the longitudinal velocity of electrons is higher by the factor $(m_i m_e)^{1/2}$ than that of ions, the angle $\Delta\psi$ is very small in comparison with unity [see (39)]. Accordingly, since the mirrors are turned through 90° relative to one another, the first term in (41) does not contribute to the overall displacement in the back and forth motion. Only the second term contributes and can be looked upon as producing a resonance with $\Delta\psi = 0$. This resonance also leads to the random radial motion of the particles (in this case, electrons), and the amplitude of the random motion is

$$\Delta r \sim \left(\frac{\zeta r}{|\partial\Delta\psi/\partial r|}\right)^{1/2}.$$

We note that if the end plugs were not at 90° to one another, the first term in (41) would contribute to δr, and the electron transport coefficients would be much higher. This type of decompensation can occur in the end traps because their mirrors are not identical. At one end, the end trap is in tandem with the central cell, and at the other end it is in tandem with the expander. Moreover, for the same reason, the ambipolar potential distribution in the end traps is asymmetric with respect to the median plane. Accordingly, the end traps provide a large contribution to the radial transport of electrons. The situation in which the rate of radial diffusion of electrons is much greater than that of ions is fully realistic. The Boltzmann distribution is then established in the electron gas not only in the direction of the lines of force, but also in the perpendicular direction, i.e., the transverse escape of electrons is prevented by the radial electric field. Under these conditions, a very high vacuum must be maintained beyond the plasma boundary, since otherwise there will be a large heat loss by electronic thermal conduction.

When electron transport is not much greater than ion transport (or is less than the latter), the electric field can again play an important part: it adjusts itself so that the overall (longitudinal and transverse) loss of electrons and ions from each tube of force is the same.

To conclude this Section, we recall that, even when resonance transport is made negligible, we still have the problem of the so-called "anomalous transport," which is due to the development of universal drift instabilities, first investigated theoretically more than 30 years ago by Rudakov and Sagdeev.[47] Experimental information on the role of these instabilities in open-ended traps is very scarce at present, and data on anomalous transport in tokamaks cannot be used directly because factors such as longitudinal currents and shear, which are typical for tokamaks, are usually absent in the case of open-ended traps. It is likely that the situation will become clearer in the course of the next two or three years, after the completion of the present-generation experiments with open-ended traps. It is expected that the rate of anomalous transport will decrease with increasing radius of the plasma column, measured in units of the ion Larmor radii $N = a/r_{Li}$. Designs of major open-ended traps are usually based on $N > 10$.

3.4. Attempts to find axially symmetric configurations

It is clear from the last Section that, if it were possible to find axially-symmetric MHD stable configurations for ambipolar traps, this would radically improve their prospects. Unfortunately, an acceptable solution of this kind has not yet emerged. The point is that all the existing experimental ambipolar systems incorporate quadrupole stabilizers that can hardly be regarded as facilitating the verification of par-

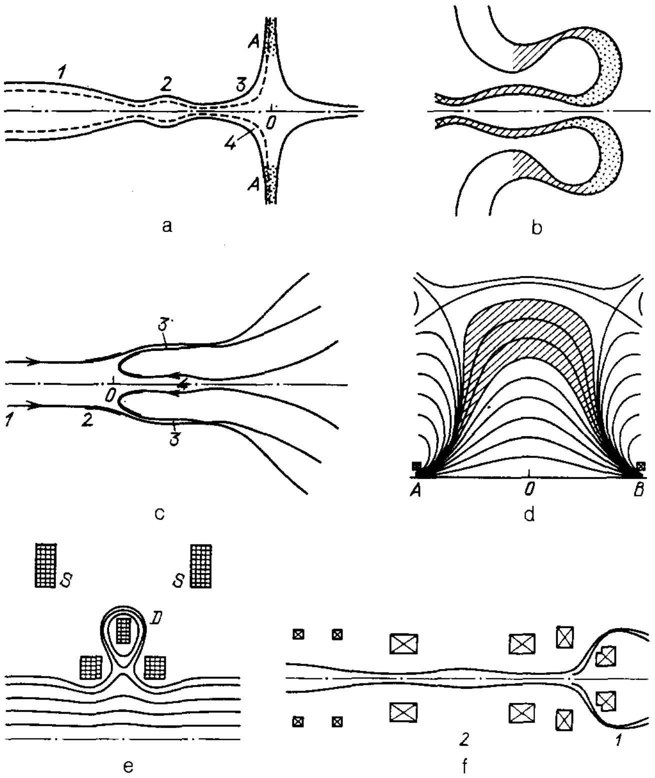

FIG. 18. Some types of axially symmetric stabilizers for ambipolar traps: a—Andreoletti-Furth stabilizer; 1—central cell, 2—plug, 3—stabilizer, 4—region occupied by connecting plasma; *A*—ring minimum of magnetic field, *O*—point of zero magnetic field; the dotted regions are occupied by the stabilizing plasma; b—the Arsenin stabilizer; shaded regions are occupied by the hot stabilizing plasma and dotted regions by the connecting plasma; double shading shows the plasma in the plug of the ambipolar trap on the left; c—a variant of the Logan stabilizer; 1—central cell, 2—stabilizing cusp region (*O*—point of zero magnetic field), 3—ring trap used to plug the ring gap in the cusp, 4—plug for axial gap; arrows show the direction of the magnetic field and thick lines show the stabilizing segments; d—example of a stabilizer in the form of a nonparaxial mirror trap; the diagram shows the magnetic configuration produced by superimposing two small mirror coils *A* and *B* and a uniform magnetic field amounting to 6% of the field produced by these coils at the point *O*; the figure shows only the upper half of the trap; e—stabilizer in the form of a "magnetic divertor"; the current in coil *D* flows in the opposite direction to the currents in the other coils; the authors of Ref. 55 ascribe the stabilization of the global interchange mode to the rapid azimuthal mixing of electrons in the neighborhood of the circle $|\mathbf{B}| = 0$; the stabilizer is placed at the midpoint of the ambipolar trap; *S*—solenoid coils; *f*—stabilizer in the form of a semi-cusp (1) and a plug (2); the central cell is located on the left.

ticular theoretical predictions for axially symmetric systems. The first tentative attempts at experimentation in this region were undertaken only very recently, and have not as yet led to clear conclusions. As far as theoretical assumptions are concerned, many of them lead to exceedingly exotic magnetic-field structures that are far removed from the "natural" mirror trap geometry (see Figs. 1b and 18), whereas other designs are based on physical principles that have not as yet been adequately demonstrated. Nevertheless, the current burst of activity leads us to hope that acceptable designs for axially-symmetric ambipolar traps will emerge in the near future.

In principle, a solution for an axially-symmetric ambipolar trap was known even in 1976, when the papers by Dimov *et al.* and Fowler and Logan were published. We have in mind the MHD trap stabilizer described by Andreoletti[48] and by Furth[49] as far back as 1963. A variant of this trap is shown in Fig. 18a (in combination with an ambipolar trap). In contrast to the usual axially-symmetric trap, in which the magnetic field has a minimum in the longitudinal direction and a maximum in the transverse direction, i.e., the center of the trap is a saddle point of the function $|\mathbf{B}(\mathbf{r})|$, the Andreoletti-Furth trap has an absolute minimum of B. The locus of such points is a circle centered on the magnetic axis, i.e., the region in which the plasma is confined is a toroid. The structure of the magnetic coils must be suitably chosen to ensure that the ring minimum appears. Of course, in the configuration illustrated in Fig. 18a, there is one further minimum, $B = 0$ at the point O, but, as already noted, here the plasma is not confined.

When the Andreoletti-Furth trap is used to stabilize an ambipolar trap, the plasma in the neighborhood of the ring minimum must be electrically connected to the plasma in the end trap. This can be done by using special "connecting" cold plasma. Since this connecting plasma is not confined in the region near the point O (and is therefore absent), the inner parts of the ambipolar trap (between the dotted lines in Fig. 18a) are not stabilized. To ensure MHD stability in these regions, a flat or radially growing pressure profile must be established in them.

Unfortunately, numerical analysis has shown that the depth of the "magnetic well" in the Andreoletti-Furth stabilizer is actually very small. There are additional difficulties (which may well be more serious) due to the fact that effective transverse dimension of the plasma in the region of the ring minimum is also small, which brings with it the possibility of microinstability development. For these reasons (and also because in 1977–1978 the difficulties with using quadrupole stabilizers were not yet fully recognized), the Andreoletti-Furth scheme was not tested experimentally and all the ambipolar traps designed in the late 1970s (and existing today) have quadrupole stabilizers.

In recognition of the difficulties associated with non-axially-symmetric systems, about twenty different designs of axially-symmetric traps have been suggested in recent years. The bibliography covering the period up to 1985 can be found in the review of Ref. 50. Some of these stabilization schemes are briefly described below.

V. V. Arsenin[51] proposed a stabilizer in which the curvature of the lines of force had an alternating sign (Fig. 18b). The advantage of this scheme is that it ensures a sufficiently deep "mean minimum B," but its disadvantage is its complicated geometry and relatively small plasma thickness.

B. Logan[52] considered a stabilizer in the form of a deformed cusp trap (Fig. 18c). To prevent plasma losses from the region near zero magnetic field (the point O in Fig. 18c), he suggested the use of ambipolar plasma confinement in a cusp with two additional traps, namely, an axial trap and a ring trap. The disadvantage of this scheme is again its spatial complexity and small plasma thickness in the confining ring plug.

A. A. Skovoroda[53] put forward an interesting suggestion in which the stabilizer was a cusp containing hot electrons and relatively cold ions. The point of using electrons as the carrier of plasma pressure is that they have a very small adiabatic region near zero magnetic field, so that their loss from the trap is relatively low.

Stabilization can also be achieved in the "natural" geometry of the magnetic mirror trap if the anisotropy of the ion distribution function is suitably chosen. In particular, by injecting fast atoms at such an angle to the magnetic field that the point at which the resulting ions come to rest lies in a region with favorable curvature of the lines of force, it is possible to ensure that a high ion pressure peak will appear at this point. In accordance with (29), the system can then be expected to be stabilized. This idea was expressed more than thirty years ago by Rosenbluth and Longmire,[27] but has remained unexploited because it was believed that it would require exceedingly high injection energy. Actually, at low injection energy, the angular distribution of the fast ions is immediately smeared out by scattering by plasma ions, and this leads to lowering of the pressure peak and to the vanishing of the stabilizing effect. The injection energy W_{inj} must therefore be chosen to be high enough, so that the fast-ion scattering time becomes much longer than their slowing-down time due to drag by electrons. This occurs when [compare (7') and (17)]

$$W_{inj} \gg T_e \left(\frac{m_i}{m_e}\right)^{1/3}, \tag{43}$$

i.e., the ion energy must indeed be very high. Pessimism in this area was deepened as a result of calculations performed by Hinton and Rosenbluth[54] who showed that, for certain particular magnetic field configuration, condition (43) had to be satisfied with an enormous (by a factor of about 50!) numerical margin.

It is only recently[55] that it has become clear that the above negative evaluation of this particular stabilization method may have to be revised. This new approach involves the optimization of the magnetic field profile and the "pumpout" of part of the scattered fast ions from a special additional magnetic trap. This can be used to reduce W_{inj} down to values of the order of $T_e(m_i/m_e)^{1/3}$. The future prospects of this stabilization technique are probably determined by the extent to which the narrow angular distribution of fast ions will respond to the development of microfluctuations.

A number of stabilization schemes have been designed to ensure MHD plasma stability only with respect to the most hazardous "global" perturbations, assuming that small-scale modes will be stabilized by the FLR effect (see Sec. 2.3).

It was suggested in Ref. 56 that the "global" mode could be stabilized by making the end mirror trap sufficiently "thick" (Fig. 18d), so that nonparaxial effects become significant [the existence of these effects is reflected already in (32) which contains a positive quantity on the right-hand side]. The important point here is that the plasma boundary does not lie too close to the axis of the system or to the separatrix. For a "stepped" pressure profile, the boundary must lie in the shaded region of Fig. 18d. Remarkably, it is not only the global mode, but also several of the first flute modes, that become stable under these conditions.

The nonparaxial stabilization effect occurs for a wide class of ion distribution functions, including the isotropic function. The absence of stringent limitations on the shape of the distribution function is an important advantage of the nonparaxial stabilizer. Its disadvantage is that the magnetic field on the equatorial plane of the trap in its outer region is lower by factors two or three than on the axis. This restricts the range of possible plasma pressures in the trap.

The divertor proposed in Ref. 57 (Fig. 18e) is outwardly similar to the nonparaxial stabilizer, and has recently been tested experimentally on the TARA system.[58] There is, however, a fundamental difference: the latter method necessarily requires that the plasma extends to the zero magnetic field and fills a certain part of the lines of force outside the separatrix. While this condition is quite readily satisfied for the moderate plasma parameters in the TARA installation, this stabilization method gives rise to difficulties for reactor plasma parameters.

G. I. Dimov and P. B. Lysyanskiĭ[59] have proposed stabilization of the global mode by a semi-cusp (Fig. 18f). The advantage of this system is the simplicity of its coils, but its disadvantage is the small plasma thickness in the region of the ring plug.

It is clear that many methods are available for stabilization of global modes (see also the review in Ref. 50). However, with all their differences, as well as advantages and disadvantages, these methods have one common feature: stabilization of the higher modes by FLR effects has not been adequately examined experimentally, and one cannot at present guarantee the absence of activity in the higher flute modes.

A unique method of stabilizing flute perturbations by high-frequency ponderomotive forces has been developed at the University of Wisconsin.[60] A stabilizing effect is observed when the h.f. field approaches the ion cyclotron frequency.

We end this Section by concluding that there are no fundamental obstacles to the development of an axially-symmetric MHD stabilizer. The question is: to what extent will a particular stabilizer fit into a real ambipolar trap without giving rise to undesirable side effects (such activization of plasma microinstabilities). In this respect, stabilization methods that preserve the "natural" ambipolar-trap geometry seem to be preferable. Real experiments are now needed.

3.5. Thermal barriers

Soon after the first experimental confirmation of the efficacy of the ambipolar method of confinement, an improvement that could be used to reduce the demands on the plasma density in the end plugs was suggested.[61] Suppose that the injection of atoms into the end plug occurs at a certain angle (usually about 45°) to the magnetic field. The ions produced as a result of the capture of atoms execute oscillations along the magnetic field and, because of the small angular spread of the original beam, the ion density has a minimum in the median plane of the plug and two maxima near the points at which the ions come to rest (Fig. 19). Ions produced in this way are commonly referred to as sloshing ions.

Next, let us suppose that a strong source of additional electron heating (e.g., an electron synchrotron resonance source) is turned on in the region of the more distant (relative to the central cell) density peak. The electron temperature T_e^* in the region of this more distant density peak can become appreciably higher than the electron temperature elsewhere in the trap. The ambipolar potential drop between the central cell and the more distant density maximum is given by (Fig. 19)[14]

$$\Delta\varphi = \frac{T_e^*}{e} \ln \frac{n_{max}}{n_{min}} - \frac{T_e}{e} \ln \frac{n_0}{n_{min}}, \tag{44}$$

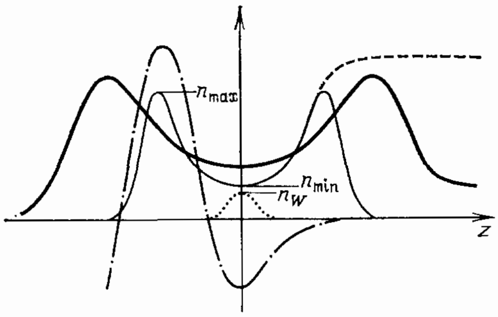

FIG. 19. Thermal barrier of an ambipolar trap. The figure shows the distribution of the magnetic field (thick line), "sloshing"-ion density (solid line), total ion density, including the "sloshing" ions and ions in the central cell (dashed line), the potential (dot-dash line), and the density of "ultrafast" electrons (dotted line).

where n_0 is the plasma density in the central cell. When T_e^* is several times greater than T_e, the second term in (43) can be neglected, and we find that the ambipolar trapping potential can be made very high even for $n_{max} < n_0$: the essential quantity is then n_{max}/n_{min} and not n_{max}/n_0 [see (37)]. It follows from the discussion given in Section 3.1 that a reduction in n_{max} is extremely desirable because it brings with it the possibility of a lower magnetic field and lower injection power.

The difficulty that arises when this scheme is implemented is due to the high rate of energy transfer between the hot electrons at temperature T_0^* and the main mass of electrons. Even when $n_{min} \sim 0.2\, n_{max}$, which is difficult to implement, the flux of hot electrons in the central cell is very high, so that the necessary difference between T_e^* and T_e can be maintained only by expending an unrealistic amount of power to heat the electrons in the region of the outer density peak. It was suggested in Ref. 61 that this difficulty could be obviated by setting up a further electron population in the region of the density minimum with characteristic energy W_e much greater than both T_e^* and the injection energy of the fast ions. Because of the high energy of these "ultrafast" electrons, their motion is not affected by the ambipolar electric field, and their longitudinal density profile is unambiguously determined by their angular distribution. By making the distribution sufficiently anisotropic $W_{e\perp} > W_{e\parallel}$, it is possible to ensure that the ultrafast electrons become concentrated near the magnetic-field minimum (see Fig.19). Let n_w denote the maximum electron density. The quasineutrality condition then shows that the density of the remaining electrons at this point must be $\Delta n = n_{min} - n_w$, and n_{min} must be replaced with Δn in (44). By choosing the density n_w to be sufficiently close to n_{min}, it is possible to ensure a significant reduction in Δn and, correspondingly, a reduction in the energy transfer between the hot and other electrons even for values of n_{min} that are not very low. As far as the "ultrafast" electrons are concerned, they have a weak interaction with the other particles because of their high energy.

Because the above experiment is based on suppression of heat transfer between two electron populations, the scheme is often referred to as the "thermal barrier" scheme.

The "ultrafast" electrons are produced by electron-cyclotron heating, using the second harmonic of the cyclotron frequency. Here it is important to note that, in the early 1970s, M. S. Ioffe *et al.*[62] used a cloud of anisotropic electrons to reduce the local plasma potential. The resulting potential well was then used to confine thermal ions stabilizing the DCI (see Section 2.4).

Since the lifetime of the main ion population confined by the electrostatic potential peaks is long in comparison with the ion scattering time τ_{ii}, collisions of "thermal" ions tend to fill the potential well in the region of the thermal barrier, to increase the ion density in this region, to increase m_{min} and, as a consequence, to cause the disappearance of the thermal barrier. It was suggested in Ref. 61 that this process can be prevented by removing the ions trapped in the potential well so that their density is always low. This is referred to as pumpout. This obviously leads to some energy losses. Roughly speaking, the power lost from one of the end traps is $V_e n_0 T_i/\tau_{ii}$ where V_e is the volume of the end trap. The power loss is acceptable when V_e is small enough.

In practice, the trapped ions can be pumped out by utilizing their charge exchange with atoms injected into the plasma with the view to producing a sloshing-ion population. Other pumpout methods must be used in reactor-scale installations in which the injection energy is high and the charge-exchange cross sections are small. Baldwin[63] has suggested that this can be achieved by exposing the plasma to a high-frequency non-axially-symmetric magnetic field of frequency chosen to be close to the frequency of longitudinal oscillations of the trapped ions, or the frequency of their drift motion around the magnetic axis. This gives rise to rapid loss of trapped ions by a mechanism analogous to that considered in the preceding Section. This pumpout method is referred to as the drift method.

The entire arsenal of ideas and experimental methods described above (with the exception of drift pumpout) has been tested on the TMX-Upgrade (Livermore) and Gamma-10 (Tsukuba) ambipolar traps, specially built to verify the efficacy of thermal barriers. It was shown that the potential distribution along the length of the device typical for thermal barriers could indeed be produced (see Fig. 19) and longitudinal ion losses from the central cell could be substantially suppressed even for $n_{max} > n_0$ (the notation is explained in Fig. 19).

On the other hand, all this was achieved only for relatively moderate plasma parameters in the central cell (density n_0 between 2×10^{12} and 3×10^{12} cm^{-3}, electron temperature $T_e \sim 100$ eV). For reasons that are not as yet fully understood, the structure of the thermal barriers disintegrated when attempts were made to increase the plasma density. Work directed toward improving the thermal barriers is continuing on the Gamma-10 installation.

There is some fear, that when thermal barriers will be used in reactor-scale installations, there will be difficulties with some oddities of the phase space, the presence of several groups of particles in it, and the presence of many peaks and valleys. This may well give rise to some unexpected complications. One of the difficulties is that the different microinstabilities will be difficult to describe theoretically in a complicated system of this kind. Another danger (noted in Ref. 64) is the possible appearance of Debye potential discontinuities, inclined at an angle to the magnetic lines of force: the reduction in the magnetic moment of ions as they pass through this type of discontinuity can produce an anomalously rapid angular spreading of the ions (see Ref. 65).

A further unfavorable effect was noted in Ref. 66, i.e., the accumulation of impurities with $Z>1$ in the region of the thermal barrier. The point is that, for ions with $Z>1$, the potential well in the region of the thermal barrier is deeper by a factor of Z than for protons. Accordingly, for the longitudinal potential distribution shown in Fig. 19, impurities from the entire installation will aggregate in this region. Since charge-exchange pumpout does not affect the impurities (the charge-exchange cross section is small), the concentration of impurities in the median plane of end traps in present-generation installations can give rise to strong scattering of "ultrafast" electrons and sloshing ions, and this will lead to the disappearance of the barriers. To avoid this effect in reactor-scale installations, it will be necessary to introduce drift pumpout of impurities.

The present author's view is that the complete implementation of the thermal barrier scheme will not be possible, but its basic principles relating to control of the longitudinal potential distribution will find applications in open-ended traps and, possibly, also in other types of fusion devices.

4. CUSPS

The attraction of the cusp trap (see Fig. 10) as a device for plasma confinement is that, since it is axially symmetric and has a relatively simple magnetic configuration, it ensures the macroscopic stability of plasma with respect to the most hazardous (flute) perturbations (see Section 2.3). Unfortunately, the presence of a null point at the center of the cusp gives rise to rapid plasma loss from lines of force passing close to this point, followed by the flute instability of the interior of the plasma. Lavrent'ev (see Ref. 67) has suggested that this could be overcome by electrostatic plugging of the ring gap and of the axial apertures of the cusp with special electrodes. The principle of this suggestion is illustrated in Fig. 20, which shows on an enlarged scale the region of the ring gap of the cusp. A high negative potential is applied to electrode 1 relative to the grounded electrode 2. When the transverse size of the gap a is small in comparison with the Debye radius,

$$a \ll r_{\rm D}\,, \qquad (45)$$

the plasma potential in the gap settles at a level close to the potential of the ground, and the plasma electrons cannot reach the plug 1 (the magnetic field in the gap is assumed to

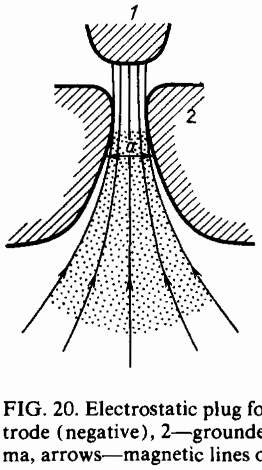

FIG. 20. Electrostatic plug for the annular gap of the cusp: 1—plug electrode (negative), 2—grounded electrode with gap. Dotted region—plasma, arrows—magnetic lines of force.

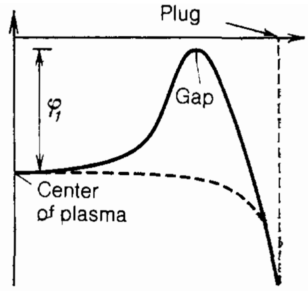

FIG. 21. Distribution of potential along a magnetic line of force emerging from the annular gap. Dashed line shows the variation in the potential due to the trapping of electrons in the potential well in the gap.

be so strong that the electrons do not strike the side walls of the gap). Axial apertures can be cut off in a similar way. The result is that the electrons are confined to the system. As far as ions are concerned, they leave the plasma until its interior becomes negatively charged, at which point ion losses cease as well (more accurately, they become exponentially small). It is clear that the particle distribution functions are then nearly Maxwellian and cannot be a source of microinstabilities.

However, this gives rise to the further problem that the potential distribution along a line of force (Fig. 21) corresponds to the presence of a deep potential well for electrons in the region of the gap. Since the electron lifetime must exceed the electron-electron collision time by orders of magnitude, this well will be filled with electrons until the electron density becomes greater by the factor $\exp(e\varphi_1/T_e)$ than the density at the center of the plasma. Condition (45) is then violated (for a realistic gap width), and the potential distribution is as shown by the dashed line in Fig. 21, i.e., the barrier for the ions disappears.

This difficulty can be overcome by continuously removing electrons from the gap. It may be that there is some natural gap-clearing mechanism such as the diocotron instability which removes trapped electrons. At any rate, experiments with a number of cusp machines have shown that the annular gap could be plugged without taking any special measures to remove trapped electrons. The experiments performed by the Ioffe group on the ATOLL installation, which will be described later, have produced direct experimental verification on the diocotron instability.[68]

The axial apertures of the cusp are more difficult to plug than the annular gap because the diameter b of the axial aperture is much greater than the width a of the annular gap, and the condition analogous to (45) is more difficult to satisfy. Actually, conservation of magnetic flux gives $\pi RaB_a = (\pi b^2/4)B_b$ where R is the radius of the annular gap and B_a and B_b are the magnetic fields in the annular gap and in the axial aperture, respectively, i.e., $b/a = 2[(R/a)(B_a/B_b)]^{1/2}$. Since, for technical reasons, the field in the axial aperture cannot exceed the field in the annular gap by orders of magnitude, it follows that indeed $b \gg a$. In actual experiments, the diameter of the axial aperture is usually much greater than the Debye radius, which means that the aperture cannot be plugged electrostatically. Nevertheless, the loss through the axial aperture is usually smaller than one would expect from simple considerations based on the nonconservation of the adiabatic invariant of ions (see Ref. 69).

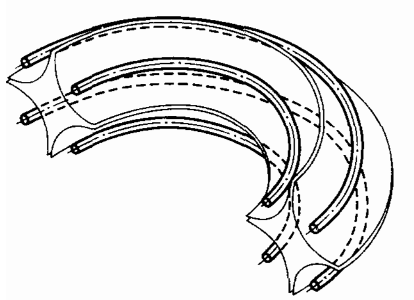

FIG. 22. Toroidal cusp. All four gaps have approximately the same width.

A possible explanation is that the motion of the particles is influenced by the electric field acting across the magnetic surfaces. The picture is still not entirely clear.

The problem of very large axial apertures can be avoided by using the toroidal cusp (Fig. 22). The ATOLL installation mentioned above is based on this principle. Experiments on this system have brought to light a further problem, namely, anomalously fast plasma loss across the magnetic field when the spatial scale of density variation is comparable with the Larmor radius of ions. This loss is due to ion scattering by short-wave fluctuations excited as a result of the relative drift of electrons and ions.[70] The ATOLL data are of general interest for all systems in which the parameter N [see (34)] is close to unity.

In view of all the above physical problems and the technical difficulties with making plugging electrodes in the case of reactor-scale plasma parameters, the cusp trap does not seem to be very promising as a reactor system. Nevertheless, cusp traps continue to attract considerable attention, mostly as stabilizing elements for ambipolar traps (see Section 3.4). Losses through the gaps are either plugged by additional mirror traps, or simply tolerated, since they can be compensated by increasing the length of the central cell.

A separate class of machines consists of cusps with high pressure plasma whose function is to exclude the magnetic field completely from the plasma-confinement region. This is closely related to work on pulsed systems. Further information can be found in the review of Ref. 71.

5. TRAP WITH ROTATING PLASMA

Plasma can be made to rotate around the mirror-trap axis by applying a transverse electric field to it. Since any magnetic surface in the equatorial plane is at a greater distance from the axis than it is in the mirror, the centrifugal force tends to "rake up" the plasma toward the equatorial plane of the trap. If the magnetic surfaces are equipotentials (the angular velocity is then constant along the magnetic surface), it is readily shown that the ion confinement condition is

$$v_{\parallel}^2 \leqslant v_{\perp}^2 (R-1) + \omega^2 a^2 \left(1 - \frac{1}{R}\right), \tag{46}$$

where $v_{\parallel}$ and $v_{\perp}$ are the components of the ion velocity in the equatorial plane of the trap (in the rotating coordinate frame), ω is the angular velocity, and a is the radius of the magnetic surface in the equatorial plane. It is clear from (46) that, for an appreciable reduction in the ion loss, the velocity ωa must exceed the thermal velocity of ions by a factor of 1.5–2. Since $\omega a = cE/B$, where E is the radial electric field and B is the magnetic field in the equatorial plane, we conclude that, under the conditions of a thermonuclear fusion reactor ($B = 3$ T, $a = 1$ m, $T = 10$ keV), the radial electric field necessary for plasma confinement must be relatively high ($E > 60$ kV/cm), and the potential drop across the plasma must exceed several million volts.

Such a large potential difference is difficult to establish in a thermonuclear fusion installation but, if this could be done, the result would be much lower longitudinal plasma loss. Moreover, the ion distribution in the rotating coordinate frame would then be nearly Maxwellian, which would mean that cone type instabilities would not appear in the plasma.

A potential drop of 400 kV across 10 cm of plasma has been produced in the PSP-2 installation[72] at the Novosibirsk Institute of Nuclear Physics (it is clear from Fig. 23 that the plasma was tubular). The electric field in the plasma was produced by a special set of electrodes (5 in Fig. 23). When the magnetic field in the equatorial plane was $B = 15$ kG, a rotational velocity of 10^8 cm/s was achieved for a typical ion energy in the rotating frame of the order of a few keV. The plasma density was about 10^{11} cm^{-3}.

Rotating plasma can be a source of different hydrodynamic instabilities, the most hazardous of which is the centrifugal instability that arises when the plasma density decreases in the radial direction. Several methods have been proposed for stabilizing this instability. We shall consider one of them[73,74] in which the density and rotational velocity profiles are such that, in the region in which rotation occurs, the density increases toward the plasma periphery, and the fall in the density near the outer boundary of the plasma occurs in a region in which there is no longer any rotation (see Fig. 24). The centrifugal force then has a stabilizing effect and improves the overall stability of the system, in-

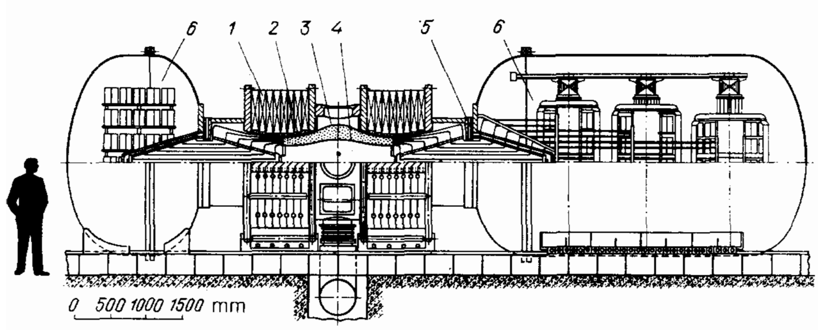

FIG. 23. Diagram of the PSP-2 installation: 1—magnetic field coils, 2—vacuum chamber, 3, 4—liners, 5—set of end electrodes used to produce the required radial potential distribution, 6—high-voltage supply system.

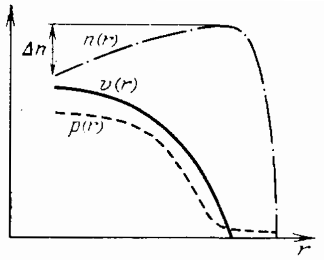

FIG. 24. Distributions of density n, rotational velocity v, and plasma pressure p ensuring the stability of internal regions with respect to flute perturbations.

cluding its stability with respect to flute perturbations. It can be shown[73] that the radial density drop necessary for suppressing flute instability can be estimated from the formula

$$\frac{\Delta n}{n} \gg \frac{v_{\mathrm{Ti}}^2}{v^2}\,\frac{a^2}{L^2} \quad (47)$$

where L is the separation between the mirrors, a is the plasma radius, and v is the characteristic rotational velocity. Since it is assumed that $v \gtrsim v_{\mathrm{Ti}}$ and $a \ll L$, a relatively small density drop is necessary to stabilize the flute instability. Of course, other methods of suppressing flute instability have to be used in the outer region of the plasma (where there is no rotation), but this situation is made easier by the fact that the plasma temperature in this region is low. Experiments performed on the SVIPP installation at the Novosibirsk Institute of Nuclear Physics[75] have confirmed that the stabilization mechanism based on (47) is valid.

Estimates based on (47) show that plasma rotation effects can be significant not only in centrifugal traps but also in all other types of open-ended trap, including the long solenoid. Actually, since a radial potential drop of the order of T/e is usually present in the plasma, the plasma rotates with velocity $v \sim cT/eBa$. Although this velocity is small in comparison with the thermal velocity of ions, the right-hand side of (47) may become less than unity for sufficiently large L. In the standard situation, the plasma density increases in the radial direction, i.e., rotation has an unfavorable effect on plasma stability. This means that control of the radial electric field, distribution, including the use of end electrodes or other similar means to produce a sharp reduction in the field in the plasma, may be required not only in the centrifugal, but also certain other types of open-ended trap. On the other hand, one is then concerned with significantly lower potential drops.

Returning now to centrifugal traps, we may conclude that, although their own thermonuclear future is not absolutely certain, mostly because the rapid rotation of the plasma is nevertheless a strong source of disequilibrium that appears in both coarse and fine instabilities, research using these machines has yielded a crop of new ideas and technical solutions that will turn out to be useful in other types of open-ended trap.

6. SYSTEMS FOR THE CONFINEMENT OF PLASMAS WITH A SHORT MEAN FREE PATH

6.1. Multimirror trap

It is clear from the discussion given in Sections 2 and 3 that the confinement of plasmas in mirror traps is associated with a significant departure of the ion distribution function (and, in the case of thermal barriers, the electron distribution function as well) from its equilibrium form in velocity space. This may become (and frequently is) a source of microfluctuations that produce fast longitudinal plasma losses, and this casts a shadow on the open-ended mirror trap as a possible reactor system.

In response to these difficulties, several open-ended systems have been considered in which L is significantly greater than the mean free path of ions λ_{ii}. The distribution of ions in the co-moving reference frame is then nearly the local Maxwellian distribution, and the problem of cone instabilities (and, generally, instabilities in velocity space) no longer arises.

The simplest example of this is a segment of a straight tube containing a uniform magnetic field (Fig. 25a), i.e., a system in which longitudinal confinement does not occur at all. Let us suppose that a blob of plasma is introduced at the initial time into the central portion of the tube (length L). Its lifetime will be of the order of its longitudinal propagation time L/v_{Ti}. For large values of L, this may be sufficient to produce values of Q exceeding unity.

In principle, the scheme shown in Fig. 25a can operate under steady state conditions if the thermonuclear plasma is introduced at one end and the reactions occur in a time L/v_{Ti} (this is the "flightotron" proposed by Morosov[76]). Unfortunately, $Q \gtrsim 1$ can be achieved for reasonable plasma densities only for very long systems (the length must exceed 30 km when, say, $N = 10^{17}\ \mathrm{cm}^{-3}$).

To reduce the rate of longitudinal plasma expansion, Budker, Mirnov, and the present author[6] suggested in 1971 that the uniform magnetic field system could be replaced by a set of connected traps (Fig. 25b) in which each individual mirror trap had a length L satisfying the condition

$$l \ll \lambda_{\mathrm{ii}}. \quad (48)$$

Particles in each trap execute finite motion between the mirrors and, when (48) is satisfied, they succeed in completing several oscillations between the mirrors between successive scattering events. Under these conditions, the transport of matter (expansion of the plasma) along the axis of the system is entirely due to drifting particles. The expansion of the plasma is accompanied by the drag on the drifting particles by the trapped particles. The latter, in their turn, transfer the momentum that they receive to the magnetic field. Conse-

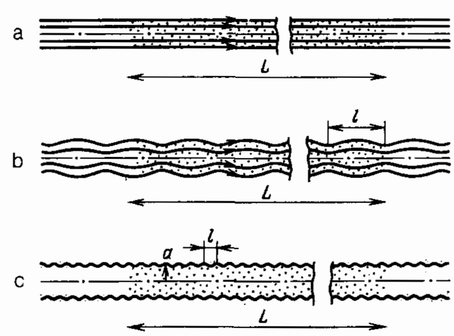

FIG. 25. Some systems for dense plasma confinement: a—straight solenoid, b—multimirror trap, c—linear system with a "rough" magnetic wall ($l \ll a$).

quently, there is a sense in which the plasma experiences drag by the magnetic field. Let us begin by considering the case of relatively small mirror ratios $R-1\sim 1$. The drag force per particle is then $F_{dr}\sim m_i u v_{Ti}/\lambda_{ii}$ where u is the drift velocity of the particles which, for $R-1\sim 1$, is of the order of the macroscopic plasma expansion velocity in the longitudinal direction. Equating nF_{dr} to the plasma pressure gradient $\partial nT/\partial z\sim n m_i v_{Ti}^2/L$, we find that the expansion velocity is $u\sim v_{Ti}\lambda_{ii}/L$. Hence it is clear that the introduction of even moderate corrugation ($R-1\sim 1$) produces an appreciable reduction in u. This is accompanied by a change in the nature of the motion itself: inertial expansion is replaced by slow diffusion, similar to the expansion of a gas in a porous medium. The longitudinal pressure gradient is balanced by the drag on the plasma by the magnetic field.

An increase in the mirror ratio produces a further reduction in expansion velocity, both as a result of the lower number of drifting particles and the shorter mean free path for scattering into the loss cone [λ_{ii} in the formula for the drag force is replaced with λ_{ii}/R; cf. (7)]. Calculations[77] show that, for "point mirrors," i.e., mirrors whose length is short in comparison with the length of an individual trap, the expansion velocity falls as $1/R^2$ for $R\gg 1$. The plasma free-expansion time is then given by

$$\tau\sim R^2\frac{L}{\lambda_{ii}}\frac{L}{v_{Ti}}. \tag{49}$$

It is clear that the lifetime increases by the factor R^2L/λ_{ii} as compared with the uniform field [to avoid misunderstanding, we recall that (49) is valid only for large mirror ratios, $R-1\gtrsim 1$, so that it cannot be taken to the limit of a uniform field].

The effect persists even for shorter mean free paths of ions[6,77] $\lambda_{ii}\lesssim l$. The drag of the plasma on the magnetic field is then due to longitudinal viscosity of the ions. The viscosity effect becomes insignificant only for very short mean free paths $\lambda_i\lesssim l^2/L$ (this estimate refers to $R-1\sim 1$). In this range of parameter values, the motion again takes the form of free expansion.

The multimirror confinement scheme for $\lambda_{ii}>1$ was discussed independently of Ref. 6 by Logan, Liebermann, Lichtenberg, and Makhijani.[7] Earlier still, a multimirror magnetic field configuration was discussed by Post,[78] but the principal effect, i.e., diffusion scaling of the plasma lifetime ($\tau\sim L^2$), was not noticed (although correct ideas were expressed in that paper on the lower sensitivity of plasmas with finite mean free path to different types of microinstability).

Among other precursors of Refs. 6, 7, and 77, we note the paper by Tuck[79] who considered longitudinal plasma confinement in a θ-pinch with zero magnetic field in the interior of the plasma, i.e., the plasma was confined radially by a "magnetic wall" (the ions traveled in straight paths between successive collisions with the wall). Tuck suggested that the magnetic wall could be made "rough" with a characteristic longitudinal irregularity smaller than the plasma radius a (see Fig. 21c). After several reflections from these irregularities, an ion changes its direction of motion along the axis of the system, i.e., the longitudinal expansion of the plasma is again diffusional in character. The advantage of this scheme is that the "step" of the random walk of the motion of the ion in the axial direction is determined by the depth and the spatial scale of the surface corrugations, and is independent of the ion mean free path, i.e., this method can also be used for $\lambda_{ii}\gg L$. Unfortunately, it is difficult to produce a plasma column with zero magnetic field in its interior and a sufficiently sharp boundary between the plasma and the magnetic field. The most hazardous factor here is the excitation of instabilities in the boundary layer between the plasma and the magnetic field (see Ref. 70). At any rate, there have been no experimental implementations of Tuck's idea.

The publication of Refs. 6, 7, and 77 was soon followed by a number of special experiments[80,82] in which the longitudinal expansion of cold alkali plasma in a multimirror magnetic field was investigated (this medium was chosen because the condition $\lambda_{ii}<L$ could be satisfied at low plasma temperature for low plasma density and moderate length L of the system). It was found that there was good agreement between the experimental results and theoretical calculations (Fig. 26). Similar experiments were subsequently run with hydrogen plasma[83] ($n\lesssim 10^{14}\,\mathrm{cm}^{-3}$, $T\sim 10\,\mathrm{eV}$), and the results were again in satisfactory agreement with calculations.

Let us now briefly consider the reactor-scale param-

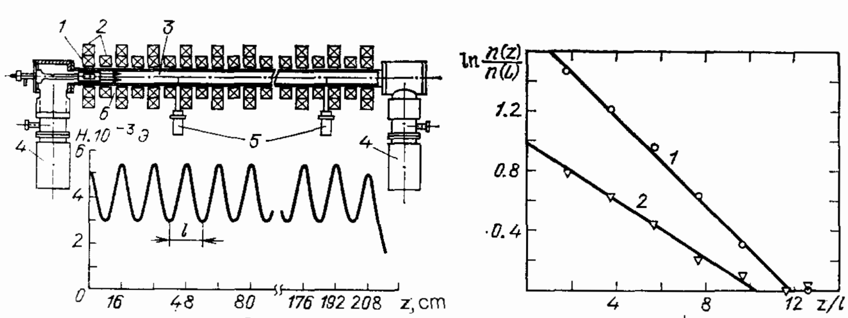

FIG. 26. Experiment on the Shchegol installation at the Novosibirsk Institute of Nuclear Physics, Siberian Division, Academy of Sciences of the USSR: a—principle of the experiment, 1—ionizer, 2—magnetic field coils, 3—vacuum chamber, 4—pumps, 5—mobile Langmuir probes, 6—device for supplying a jet of cesium atoms to the ionizer; the magnetic field distribution along the length of the installation is shown under the apparatus; the plasma is produced at the hot tungsten plate 1 and flows along the multimirror magnetic field toward the cold absorber at the other end of the installation; b—steady-state distribution of plasma density along the axis of the system; in accordance with theoretical predictions, the exponential function $n(z)$ is realized when the ion mean free path is short enough. Free molecular flow occurs in the end section, and the density is independent of z; curve 2 differs from curve 1 by the lower plasma current from the ionizer (see Ref. 81).

eters of multimirror traps. If we take the traditional value $n \sim 10^{14}$ cm^{-3} for a quasistationary thermonuclear installation, then the condition $\lambda_{ii} < L$ demands that the installation must have an unrealistic length, namely, $L > 3 \times 10^4$ m. It is therefore desirable to increase the plasma density. For quasistationary thermonuclear systems, the plasma density is restricted by the condition $2nT \lesssim B^2/8\pi$ where B is the magnetic flux density that confines the plasma, preventing it from expanding in the radial direction. When $B \sim 15$ T this condition restricts the plasma density to $n \sim 10^{16}$ cm^{-3}. Several American publications (see for example Ref. 84) have discussed multimirror reactors with plasma of this density. It was assumed that steady-state plasma was sustained by the injection of neutrals; the heating of plasma by alpha particles was taken into account. The general conclusion is that the length of the multimirror reactor with $Q = 5$ will be about 1 km under stationary conditions.

A different approach is being developed in the Soviet Union. Here, the length of the installation is reduced by using still denser plasma[6]: $n \sim 10^{17}$ cm^{-3}. Since magnetic confinement of this plasma by realistic magnetic fields is not possible, it was suggested[6] that the old idea of "wall confinement" should be employed, in which equilibrium in the radial direction is assured by direct mechanical contact between the plasma and the walls of the working chamber, and the magnetic field is used only to suppress thermal conduction to the chamber walls, i.e., $\beta \gg 1$. Of course, this will have to be a pulsed reactor because of the considerable mechanical and thermal load on the chamber wall. It was proposed that a high-current electron beam injected at one end of the installation would be used for the rapid heating of the plasma. An interesting property of the wall confinement installation is that the MHD stability of the plasma is assured even in an axially-symmetric system (because of the direct mechanical contact between plasma and the walls).

The importance of this last approach to open-ended traps is that it is based on a totally different (nontraditional) technique which promises to be cheaper. Moreover, it enables us to extend the ideas and methods associated with open-ended traps to a totally different range of plasma parameters (such as β and $\omega_{Bi}\tau_{ii}$) in which the problem of anomalous transverse transport (see Section 3.3) will be favorably resolved. Calculations on the pulsed thermonuclear reactor with multimirror confinement can be found, for example, in Ref. 85.

The entire range of phenomena that occur in the case of wall confinement of dense plasmas will be studied on the GOL-3 installation that is being built at the Novosibirsk Institute of Nuclear Physics.[86] Its parameters will be as follows: length 22.5 m, plasma diameter 26 cm, average magnetic field over the length of the system 6 T. The expected plasma parameters are: density 10^{17} cm^{-3}, temperature 1 keV, lifetime 100 μs.

The characteristics of the multimirror reactor (both pulsed and continuously operating) can be improved by introducing a certain amount of impurities[87] with $Z \gg 1$. Of course, this gives rise to an increase in bremsstrahlung from the plasma, but there is a simultaneous reduction in the mean free path and, according to (49), an increase in confinement time. The introduction of the optimum amount of impurity can result in an increase in the energy multiplication factor Q by a factor of 2–3.

6.2. Gas-dynamic trap

A further example of the open-ended trap for the confinement of plasmas with a short mean free path is the so-called gas-dynamic trap (GDT) proposed in Ref. 88. The GDT is a mirror trap with a very large mirror ratio and length L satisfying the condition

$$L \gg \lambda_{ii} \frac{\ln R}{R} . \tag{50}$$

This condition signifies that the length of the system is greater than the length for ion scattering through an angle equal to that of the loss cone (the factor ln R is due to small-angle Coulomb scattering). Accordingly, the ion distribution function is nearly Maxwellian at all points, with a possible exception of the region near the "throat" of the mirror, and the plasma lifetime is defined by analogy with the time taken for gas to issue from a vessel with a small aperture (i.e., as the ratio of the total number of particles nV in the vessel of volume V to the gas current $nv_T S$ through an aperture of area S). In the present case, this gives

$$\tau \sim \frac{RL}{v_{Ti}} . \tag{51}$$

The pure gas-dynamic origin of (51) is the reason for the name of the trap.

As should be the case in the gas-dynamic situation, the collision frequency does not appear in (51) which, in a certain sense, gives the limiting plasma "consumption": this "consumption" would still not be exceeded even if microinstabilities were to appear in the plasma and the collision frequency became higher than the Coulomb value. It is also remarkable that the lifetime increases linearly with the mirror ratio (and not logarithmically as would be the case in the ordinary short mirror trap). Accordingly, it is sensible to increase the GDT mirror ratio to the maximum possible attainable value. In this sense, the possibility, mentioned in Ref. 88, of producing an MHD stable GDT with an axially symmetric configuration is very important (with the axially symmetric mirror coil, the magnetic field at the conductor is almost the same as on the axis of the system, whereas for practical designs of the quadrupole mirrors, the field at the conductor is a few times higher than on the axis).

The stabilization mechanism noted in Ref. 88 is based on the fact that the plasma density (and pressure) in the mirror and immediately behind the mirror is the same as in the central part of the trap,[15)] so that the region outside the mirror provides an appreciable contribution to the stability integral (29). The lines of force in the region outside the mirror have a favorable curvature, so that this contribution is a stabilizing one. By ensuring that the magnetic field after the "throat" of the mirror diverges rapidly, i.e., by increas-

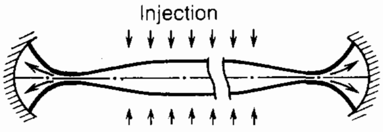

FIG. 27. Gas-dynamic trap: short arrows show the direction of neutral atoms and long arrows indicate the plasma current into the expander. The figure also shows segments of lines of force providing a favorable contribution to the "stability integral" (29).

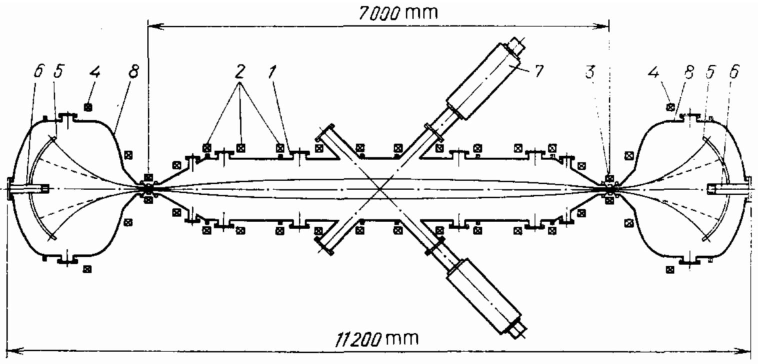

FIG. 28. Principle of the gas-dynamic trap: 1—central vacuum chamber, 2—coils of the solenoidal part of the magnetic system, 3—internal mirror coils producing 11 T, 4—coils defining the magnetic field geometry in the expanders, 5—plasma detectors, 6—plasma gun, 7—atomic beam injectors, 8—vacuum chambers of the expanders.

ing the curvature of the lines of force and, at the same time, making the transition from the uniform field to the mirror very smooth (as shown in Fig. 27), it is possible to achieve a favorable sign for the integral (29) as a whole. When this integral is evaluated for flowing plasma, we have to put $p_{\parallel} = p + \rho v^2$, $p_{\perp} = p$ where p is the gas-kinetic plasma pressure (which is isotropic in the rest frame for plasma with high collision frequency) and v is the flow velocity.

The fact that the MHD stability of the GDT can be assured by suitably choosing the shape of the lines of force in the region after the mirror was verified experimentally in 1986 by G. D. Roslyakov *et al.*, at the Novosibirsk Institute of Nuclear Physics.[89] The experiment is shown schematically in Fig. 28. The special coils (4 in Fig. 28) can be used to adjust the magnetic-field configuration in this region between the destabilizing (broken lines in Fig. 28) and stabilizing (solid lines) configurations. The magnetic field inside the trap is practically unaffected by this. In the unfavorable field distribution, strong flute-type oscillations occur in the plasma, and losses are mostly across the field. In the favorable field distribution, the energy of the flute oscillations is reduced by a factor of several tens, and the plasma losses occur through the mirrors, so that the lifetime is given by (51) (Fig. 29).

The linear dependence of τ on R was also verified in these experiments up to $R = 25$. This was extended up to $R = 74$ in the experiments reported in Ref. 90.

The GDT reactor is designed to operate in the steady-state regime. Plasma and heat losses through the mirrors will be compensated either by the injection of atomic beams (mass and energy supply) or by periodically introducing frozen pellets of the DT fuel into the plasma (energy balance can then be assured by, for example, ion-cyclotron heating). On the whole, this system is attractive both in its simplicity

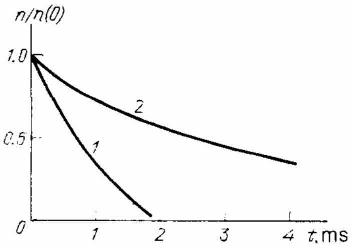

FIG. 29. Plasma density in the trap as a function of time: 1—unfavorable field distribution in the expanders (dashed lines in Fig. 28), 2—favorable curvature (solid lines in Fig. 28).

and the reliability of the confinement scheme (for example, longitudinal losses are insensitive to the development of plasma microfluctuations). Unfortunately, this simplicity has to be paid for by the fact that the GDT reactor is relatively long. If, following Ref. 91, we take the vacuum field in the mirrors to be $B_{max} = 45$ T and the field over the uniform region $B_{min} = 1.5$ T, then the typical length of a reactor with $Q = 5$ turns out to be 2–3 km. This is too long according to current ideas. However, we recall that practically the entire length of this reactor is occupied by a simple uniform solenoid producing a weak field, so that, in the final analysis, economic (and not emotional) arguments will be decisive.

The characteristics of the GDT reactor depend significantly on the limiting magnetic field that can be established in the mirrors. The relatively imposing figure of 45 T was assumed for this field in the above discussion. According to current ideas, this field can only be produced by a "warm" coil (perhaps, with a superconducting outer part), so that a certain amount of power has to be expended in maintaining the field in the mirrors. However, the field need only be produced in a relatively small volume (diameter 10–15 cm, length 20–30 cm), so that the power consumption is moderate.[91] The thermonuclear prospects of the gas-dynamic trap would become quite good if advances in the technology available for producing ultrastrong magnetic fields were to result in the generation of 60–80 T with continuously operating coils.

The length of the installation could also be reduced at the cost of some complication of the design, namely, by adding a further mirror trap at each end, with the trap length satisfying (50), but short in comparison with λ_{ii}. The form of the plasma balance equations for the additional traps (see Ref. 92) shows that the longitudinal plasma loss decreases by a factor of about two as compared with the single gas-dynamic trap. This improvement is achieved at the cost of a more complicated design and increased power consumption in the mirror coils.

7. THE OPEN-ENDED TRAP AS A NEUTRON GENERATOR

A source of neutrons based on the DT reaction[16)] would be very useful for testing materials and the constructional elements of a future thermonuclear fusion reactor. The neutron flux density produced by the neutron generator in the test zone would have to be several times greater than on the "first wall" of the reactor, where it is expected to be 0.5×10^{14}–1.5×10^{14} cm^{-2} s^{-1} (Ref. 93). Of course, in addition to neutrons from the DT reaction, there will also be secondary neutrons whose spectrum and flux will be determined by the plasma environment. The figures we have just given refer to the primary neutron flux.

The neutron source need not necessarily have a positive energy yield, i.e., $Q > 1$. The important factors are the neutron flux density in the test zone and engineering simplicity.

The attractiveness of the open-ended trap as a neutron source is based on the following considerations. The number S of neutrons produced per unit plasma volume per unit time for a given shape of the energy distribution of deuterons and tritons is proportional to n^2:

$$S \propto n^2. \tag{52}$$

On the other hand, for a given energy distribution, the plasma pressure is proportional to n. Since the equilibrium conditions show that $p \sim \beta B^2/8\pi$, where B is the confining magnetic field, we find from (52) that

$$S \propto \beta^2 B^4. \tag{53}$$

It is clear that S increases rapidly with increasing β, and it is well-known that, open-ended traps have record values of β: the value $\beta \simeq 1$ has been attained in such traps (under steady-state conditions!; see Section 2.3). The neutron number S varies even more rapidly with B, and here again the axially symmetric open-ended trap has no rivals: as already noted, the field on the axis can be made very close to the field "at the conductor." Thus, open-ended traps, and especially axially symmetric traps, are indeed very promising as neutron generators.[17)]

It is therefore not surprising that there have been several recent suggestions for neutron sources of this type.[94–97] To illustrate existing possibilities, let us describe a neutron source based on the gas-dynamic trap[96] (see also Ref. 92). This source is based on the so-called "target" approach: it is assumed that the trap is filled with relatively cold ($T \sim 1$ keV) and dense ($n \sim 5 \times 10^{14}$ cm^{-3}) deuterium plasma in which fast tritons are injected with an energy of 250 keV. The tritons collide with the deuterons in the target plasma, and create neutrons and α particles. The outflow of the target plasma through the mirrors can be compensated by injecting pellets of deuterium into the system, so that it will operate in a steady-state regime.

The tritons are injected at a small angle ($\theta_* \approx 20°$) to the magnetic field, and oscillate between the points at which they come to rest. As they do so, they are slowed down by plasma electrons and are scattered by deuterons. The points at which fast tritons come to rest in the magnetic field are given by $B_* = B_0/\sin^2\theta_*$, where B_0 is the field over the uniform portion. Because the electron temperature is not high, the slowing down of the tritons by electrons occurs much more rapidly than scattering by deuterons, i.e., the width of the angular distribution of the tritons is small in comparison with θ_*. Therefore, their density has a sharp peak near the points at which they come to rest. The maximum density is limited only by the condition that the magnetic pressure must be greater than the triton pressure. As a result, the stopping point and its neighborhood become a powerful source of neutrons.

The calculations reported in Ref. 92 and 96 show that, for a total power consumption of 50–100 MW, this type of neutron source can generate a flux of 2×10^{14}–3×10^{14} cm^{-2} s^{-1} on the surface of a cylinder 15 cm in diameter and 1–1.5 m long. The separation between the GDT mirrors can be made ~ 10 m, so that the source as a whole is relatively compact. It is also important that the basic dimensionless plasma parameters [see (34)] at the stopping points of the fast tritons are not very different from those in the 2XIIB installation, and this provides additional reliability for the entire design.

The properties of open-ended traps that enable us to use them as neutron sources are also favorable in relation to the possible use of "nontraditional" thermonuclear fuel, i.e., in the first instance a mixture of deuterium and helium-3. The reaction $D + He^3 \to He^4 + p + 18.6$ is interesting in that it does not produce neutrons, and the reactor itself is virtually free from radioactivity. The word "virtually" means that if we use the $D - He^3$ mixture in equal proportions, the "side" reaction $D + D \to T + n$ will also take place, but the cross section for it is relatively small (at the temperatures at which the $D - He^3$ reaction rate is already considerable), and the number of neutrons will also be small.

The difficulty with producing a reactor based on the $D - He^3$ mixture is that the reaction has an appreciable rate only at temperatures in the range $\gtrsim 100$ keV, whereas the DT reaction proceeds satisfactorily even at temperatures of 5–10 keV. Accordingly, for a given plasma pressure, the density of the $D - He^3$ mixture will be lower by a factor of 10–20 than the density of the DT mixture, and the power released per unit volume will be very small. Hence the $D - He^3$ reactor will depend critically on the extent to which a high plasma pressure will be allowed by the confinement system. As already noted, open-ended traps seem very satisfactory from this point of view.

8. CONCLUSION

As we consider all the open-ended traps described in this review, we gain the paradoxical impression that the best of them is the simplest mirror trap, especially in its axially symmetric variant. It implements in full measure all the advantages of the open-ended trap, namely, the tempting simplicity of construction, the steady-state regime, the attainability of values of β close to unity, and the availability of a natural channel for removing impurities and thermonuclear products. Unfortunately, all these remarkable properties of the mirror trap have to be balanced against a plasma lifetime that is not long enough in relation to losses through the mirrors, and energy multiplication factor Q that cannot exceed values of the order of 1.5.

Nevertheless, these first impressions are to some extent correct: improvements of the mirror trap directed towards increasing Q have usually resulted in the appearance of additional components and, of course, in a more complicated system. Not surprisingly, this has also been accompanied by the loss of one or two of its original attractive properties. However, the last thirty years of research into open-ended traps have not been wasted: in many cases, the debits are not too significant and have been balanced by gains in Q. The most obvious example is the gas-dynamic trap which, externally, looks like a simple mirror trap (except that it is much longer) and retains all its above advantages. The only disadvantage, according to current ideas, is its length. Of course, in the longer term, this length and the thermonuclear power yield may come to be regarded as acceptable. At present, however, the system seems too long. It may therefore become necessary to complicate the system somewhat by adding two additional mirrors (see Section 6.2) and this makes it possible to reduce the length. A significant improvement in the parameters of the GDT reactor may also result from further advances in the technology used to produce ultrastrong magnetic fields, which should lead to stronger "point" plugs. In the shorter term, the gas-dynamic trap has a good chance of becoming a high-flux source of thermonuclear neutrons.

Ambipolar traps, in their axially symmetric variant, also have the basic advantages of the mirror trap except, possibly, the fact that there are again difficulties with the removal of heavy impurities (which now cannot leave the central cell because of the presence of potential barriers at the ends). If an acceptable axially symmetric configuration can be found for the ambipolar trap, then the only remaining serious problem for these systems will be stabilization of "velocity-space instabilities" in the end traps. Since, as the distribution function and the velocity space become more complicated, the range of these instabilities usually expands, attempts to improve the properties of the end mirror traps by establishing thermal barriers in them give rise to certain dangers. The original ambipolar trap (without thermal barriers) has the best chance, at least in the short term. Of course, the corresponding reactor will be relatively long.

The multimirror system with wall confinement of dense plasma loses the advantage of steady-state operation as compared with the mirror trap, and is also constructionally more complicated, but it retains the other advantages of the open-ended traps. Based on its physics this system is quite simple and reliable. Its future on the reactor scale will depend on advances in the development of materials for the first wall which will have to withstand the high pressure and the severe heat transfer conditions.

We should not lose sight of the other types of open ended trap. However, we have emphasized three systems on the basis of one important criterion: they are already represented by relatively large working (or nearing completion) experimental installations. Moreover, there are no exceptions to the peculiar continuity principle according to which machines belonging to successive generations cannot differ (in their parameters) by a factor of more than 2–3, since otherwise they become impractical. Since the current-generation open-ended trap is a very large and expensive system (of the scale of the T-15 tokamak), the choice of the design of this installation is necessarily limited to the above three varieties of open-ended trap, which will enable us to verify the basic principles on a sufficiently large scale. Of course, the other open-ended traps can contribute to this databank of ideas, which will be used to develop installations of the next generation. It is also possible that the eventual installation will be a "hybrid" and will incorporate features of different open-ended traps. However, it is unlikely that it will be the result of an extrapolation of results obtained with bench-top systems.

The next five or six years will be very important for open-ended traps. During these years at least one of the traps that already exists or is being built will prove itself as a prototype for the next generation of machines in which plasma with near-reactor parameters will be produced. Of course, one could continue to extend the range of existing installations, developing new systems of the same scale. However, it seems that enough time has been spent on applying the method of trial and error to open-ended traps: if the necessary level of confidence is not reached in the next five or six years, then the view of the present author is that interest in open-ended traps will rapidly decline, and they will leave the scene (although definitely not forever: the open ended trap is irreplaceable as a system for burning "neutron-free" fuel, but this will involve a completely different level of technology).

As far as applications of open ended traps in high-flux neutron sources are concerned, here the future is quite clear: such neutron sources can be built if required.

[1)]In particular, all development work on open-ended traps came to an end in Western Europe where there was a strong tradition of research in this area.

[2)]Strictly speaking, this belongs to a class of closed traps, originally based on the mirror trap.

[3)]This has remained a record set of parameter values to this day (including tokamaks). To be fair, however, we note that the ion lifetime was only 10^{-3} s.

[4)]We shall assume throughout this paper that this condition is satisfied and will therefore discard the subscripts e and i on electron and ion densities (we shall use the common symbol n for both).

[5)]In the case of charge exchange, the number of plasma particles remains unaltered because a plasma ion is simply replaced by a beam ion.

[6)]We use the phrase "plasma pressure" since $p_\parallel \approx p_\perp$.

[7)]By definition, these have a frequency that is low in comparison with ω_{Bi}, and a spatial scale exceeding r_{Li}.

[8)]The magnetic field configuration used in these experiments was somewhat different from that shown in Fig. 9.

[9)]For the sake of simplicity, we ignore the charge exchange trapping of beam atoms (in the course of which "thermal" atoms escape from the plasma).

[10)]The fact that Q is small also indicates that the contribution of α particles to plasma (electron) heating has no significant effect on the energy balance, since the α particles account for only 20% of P_f.

[11)]Detailed calculations of the plasma particle lifetime with electrostatic potential barriers present in the mirrors have been carried out by Pastukhov (cf. his review paper[34]). The formula for the particle lifetime for $e\Delta\varphi \gg T_i$ is referred to as the Pastukhov formula.

[12)]The function U is determined by (31) in which, in this case, the integral must be evaluated between the peaks of the ambipolar potential.

[13)]There is also a certain azimuthal displacement that does not remove the ion from the magnetic surface, but this can usually be neglected in comparison with the corresponding shift (38) on the uniform segment.

[14)]Formula (44) is approximate because the two electron distributions overlap near the density minimum, and the concept of a temperature is no longer meaningful. When $n_{min} \ll n_0, n_{max}$ the error introduced by this formula is small.

[15)]This is in contrast to the traditional mirror trap in which the plasma density beyond the mirror is vanishingly small with respect to the parameter L/λ_{ii}.

[16)]I.e., neutrons with energy of 14 MeV.

[17)]Strictly speaking, (53) refers to the case where the mean deuteron and triton energies are roughly the same. However, in the case of the "target" approach discussed below, the dependence of S on β and B remains strong.

[1]G. I. Budker, *Plasma Physics and the Problem of Controlled Thermonuclear Reactions* [in Russian], Academy of Sciences of the USSR, Moscow 3, 3 (1958).

[2]A. S. Bishop, Project Sherwood, Addison-Wesley, Reading, MA, 1958 [Russ. transl., Gosatomizdat, M., 1960].

[3]Yu. V. Gott, M. S. Ioffe, and V. G. Tel'kovskii (in Russian), Nucl. Fusion Suppl. 3, 1045 (1962).

[4]E. K. Zavoiskiĭ, S. L. Nedoseev, L. I. Rudakov, V. D. Rusanov, V. A. Skoryupin, and S. D. Fanchenko (in Russian), Proc. Third Intern. Conf. on Plasma Physics and Controlled Nuclear Fusion Research, IAEA, Vienna 2, 679 (1969).

[5]V. V. Arsenin, V. A. Zhil'tsov, and V. A. Chuyanov, *ibid*., p. 515.

[6]G. I. Budker, V. V. Mirnov, and D. D. Ryutov, Pis'ma Zh. Eksp. Teor. Fiz. **14**, 320 (1971) [JETP Lett. **14**, 212 (1971)].

[7]B. G. Logan, M. A. Liebermann, A. J. Lichtenberg, and A. Makhijani, Phys. Rev. Lett. **28**, 144 (1972).

[8]S. G. Konstantinov, O. K. Myskin, A. F. Sorokin, and F. A. Tsel'nik, Zh. Tekh. Fiz. **41**, 2527 (1971) [Sov. Phys. Tech. Phys. **16**, 2006 (1972)].

[9]T. K. Fowler and R. F. Post [Russ. transl., Fiz. Plazmy 3, 1408 (1977)] Sov. J. Plasma Phys. 3, 787 (1977).
[10]F. H. Coensgen, J. F. Clauser, D. L. Correll, W. F. Cummins, C. Gormezano, B. G. Logan, A. W. Molvik, W. E. Nexsen, T. S. Simonen, B. W. Stallard, and W. C. Turner, Proc. Sixth. Intern. Conf. on Plasma Physics and Controlled Nuclear Fusion Research, IAEA, Vienna 3, 135 (1977).
[11]G. I. Dimov, V. V. Zakaidakov, and M. E. Kishenevskii, Fiz. Plazmy 2, 597 (1976) [Sov. J. Plasma Phys. 2, 326 (1976)].
[12]T. K. Fowler and B. G. Logan, Comments Plasma Phys. Controll. Fusion 2, 167 (1977).
[13]D. D. Ryutov and D. V. Stupakov, Pis'ma Zh. Eksp. Teor. Fiz. 26, 186 (1977) [JETP Lett. 26, 174 (1977)].
[14]V. V. Mirnov and D. D. Ryutov, Pis'ma Zh. Tekh. Fiz. 5, 678 (1979) [Sov. Tech. Phys. Lett. 5, 279 (1979)].
[15]D. D. Ryutov, Nucl. Fusion 20, 1068 (1980).
[16]V. A. Chuyanov, Itogi Nauki i Tekhniki, Ser. Fizika Plasmy, VINITI, Moscow 1, 119 (1980).
[17]T. C. Simonen, Proc. IEEE 69, 935 (1981).
[18]E. E.Yushmanov, Nucl. Fusion 25, 1217 (1985).
[19]R. F. Post, *ibid.* 27, 1579 (1987).
[20]B. V. Chirikov, in: Voprosy Teorii Plasmy, Energoatomizdat, M., 1983, No. 13, p. 3 [Reviews of Plasma Physics, Consultants Bureau, N.Y., 1987, Vol. 13, p.1].
[21]B. A.Trubnikov, in: *ibid.*, No. 1, 98 (1963) [Reviews of Plasma Physics, Consultants Bureau, N. Y., 1963, Vol. 1, p. 105].
[22]N. N. Semashko, see Ref. 16, page 232.
[23]F. H. Coensgen and T. C. Simonen, *Physics of Plasma Close to Thermonuclear Conditions*, Proc. Intern. School of Plasma Physics, Varenna, Italy, 1979, CEC, Bruxelles 2, 659 (1980).
[24]S. I. Braginskii, see Ref. 21, in: No. 1, 183 [Reviews of Plasma Physics, Consultants Bureau, N.Y., 1963, Vol. 1, p. 205].
[25]L. D. Landau and E. M. Lifshitz, *The Classical Theory of Fields*, Pergamon Press, Oxford, 3rd. ed., 1971 [Russ. original, Nauka, M., 1967, Section 33].
[26]A. B. Mikhailovskiĭ, *Theory of Plasma Instabilities*, Consultants Bureau, N.Y., 1974 [Russ. original, Atomizdat, M., 1971, Vol. 2].
[27]M. N. Rosenbluth and C. L. Longmire, Ann. Phys. (N.Y.) 1, 120 (1957).
[28]B. B. Kadomtsev, *Plasma Physics and the Problem of Controlled Thermonuclear Reactions* [in Russian], Academy of Sciences of the USSR, Moscow 4, 16 (1958).
[29]M. N. Rosenbluth, N. A. Krall, and N. Rostoker, Nucl. Fusion Suppl. 1, 143 (1962).
[30]R. F. Post and M. N. Rosenbluth, Phys. Fluids 9, 730 (1966).
[31]A. B. Mikhailovskii, Nucl. Fusion 5, 125 (1965).
[32]D. Baldwin, Rev. Mod. Phys. 49, 317 (1977).
[33]B. I. Davydov, see Ref. 28 1, 77.
[34]V. P. Pastukhov, see Ref. 20, page 160 [p. 203 in Rev. Plasma Phys. Vol. 13, 1987].
[35]G. G. Kelley, Plasma Phys. 9, 503 (1967).
[36]S. Miyoshi, K. Yatsu, T. Kawabe, K. Ishii, A. Itakura, H. Izhizuka, and S. Hagiwara, Proc. Seventh Intern. Conf. on Plasma Physics and Controlled Nuclear Fusion Research, IAEA, Vienna 2, 437 (1979).
[37]F. H. Coensgen *et al.*, Phys. Rev. Lett. 44, 1132 (1980).
[38]D. D. Ryutov and G. V. Stupakov, Dokl. Akad. Nauk. SSSR 240, 1086 (1978) [Sov. Phys. Dokl. 23, 412 (1978)].
[39]A. I. Morozov and L. D. Solov'ev, see Ref. 21, 1963, No. 2, 177 [in Rev. Plasma Phys. Vol. 2, 1966].
[40]D. D. Ryutov and G. V. Stupakov, see Ref. 20, page 74 [p. 93 in Rev. Plasma Phys. Vol. 13, 1987].
[41]A. A. Galeev and R. Z. Sagdeev, see Ref. 20, No. 7, 205 [p. 1 in Rev. Plasma Phys. Vol. 7, 1973].
[42]T. C. Simonen, S. L. Allen, and D. E. Baldwin *et al.*, Proc. Eleventh Intern. Conf. on Plasma Physics and Controlled Nuclear Fusion Research, IAEA, Vienna 2, 231 (1987).
[43]T. Cho, M. Ichimura, M. Inutake, K. Ishii, A. Itakura, I. Katanuma, Y. Kiwamoto, A. Mase, S. Miyoshi, Y. Nakashima, T. Saito, K. Sawada, D.Tsubouchi, N. Yamaguchi, and K. Yatsu, *ibid.*, page 243.
[44]D. Baldwin and L. D. Pearlstein, Memorandum MFE/TC/78-189, Lawrence Livermore Lab., 1978.
[45]G. V. Stupakov, Fiz. Plasmy 5, 958 (1979) [Sov. J. Plasma Phys. 5, 534 (1979)].
[46]D. A. Panov, Preprint IAE-3535/6 Moscow, 1981.
[47]L. I. Rudakov and R. Z. Sagdeev, Dokl. Akad. Nauk. SSSR 138, 581 (1961) [Sov. Phys. Dokl. 6, 415 (1961)].
[48]J. Andreoletti, C. R. Acad. Sci. 257, 1235 (1963).
[49]H. Furth, Phys. Rev. Lett. 11, 308 (1963).
[50]D. D. Ryutov, Plasma Phys. Controll. Fusion. 28, 191 (1986).
[51]V. V. Arsenin, Pis'ma Zh. Eksp. Teor. Fiz. 37, 534 (1983) [JETP Lett. 37, 637 (1983)].
[52]B. G. Logan, Comments Plasma Phys. Controll. Fusion 6, 199 (1981).
[53]A. A. Skovoroda, Fiz. Plasmy 11, 1319 (1985) [Sov. J. Plasma Phys. 11, 755 (1985)].
[54]F. Hinton and M. N. Rosenbluth, Nucl. Fusion 22, 1547 (1982).
[55]I. A. Kotel'nikov, G. V. Roslyakov, and D. D. Ryutov, Fiz. Plasmy 13, 403 (1987) [Sov. J. Plasma Phys. 13, 227 (1987)].
[56]D. D. Ryutov and G. V. Stupakov, Pis'ma Zh. Eksp. Teor. Fiz. 26, 186 (1985) [JETP Lett. 26, 174 (1977)].
[57]B. Lane, R. S. Post, and J. Kesner, Nucl. Fusion 27, 277 (1987).
[58]R. S. Post *et al.*, Proc. Eleventh Intern. Conf. on Plasma Physics and Controlled Nuclear Fusion Research, IAEA, Vienna 2, 251 (1987).
[59]G. I. Dimov and P. B. Lysyanskiĭ, Preprint IYaF-SO AN SSSR, No. 86–102, Novosibirsk, 1986.
[60]N. Hershkowitz, D. A. Breun, D. A. Brouchous, J. D. Callen, C. Chan, J. R. Conrad, J. F. Ferron, S. N. Golovato, R. Goulding, S. Horne, S. Kidwell, B. Nelson, H. Pershing, J. Pew, S. Ross, G. Severn, and D. Sing, Proc. Tenth Intern. Conf. on Plasma Physics and Controlled Nuclear Fusion Research, IAEA, Vienna 2, 265 (1985).
[61]D. E. Baldwin and B. G. Logan, Phys. Rev. Lett. 43, 1318 (1979).
[62]M. S. Ioffe, B. I. Kanaev, V. P. Pastukhov, and E. V. Yushmanov, Zh. Eksp. Teor. Fiz. 67, 2149 (1974) [Sov. Phys. JETP 40, 1064 (1974)].
[63]D. Baldwin, *Mirror Based and Field-Reversed Approaches to Magnetic Fusion*, Proc. Intern. School on Plasma Physics, Italy, Varenna 1, 109 (1983).
[64]L. S. Pekker, Preprint IYaF SO AN SSSR, No. 80–161, Novosibirsk, 1980; Fiz. Plasmy 10, 61 (1984) [Sov. J. Plasma Phys. 10, 33 (1984)].
[65]I. A. Kotel'nikov, *Ibid.* 12, 63 (1986) [12, 34 (1986)].
[66]D. D. Ryutov, *Ibid.* 13, 1286 (1987) [13, 741 (1987)].
[67]O. A. Lavrent'ev, *Magnetic Traps* [in Russian], Naukova Dumka, Kiev 3, 77 (1968).
[68]V. V. Piterskii, V. P. Pastukhov, and E. E. Yushmanov, Fiz. Plasmy 13, 51 (1987) [Sov. J. Plasma Phys. 13, 29 (1987)].
[69]Yu. A. Azovskii, V. I. Karpukhin, O. A. Lavrent'ev, V. A. Maslov, M. N. Novikov, and M. G. Nozdrachov, *Ibid.* 6, 296 (1980) [*sic*].
[70]M. S. Ioffe, B. I. Kanaev, V. P. Pastukhov, V. V. Piterskii, and E. E. Yushamov, *ibid.* 10, 464 (1964) [10, 268 (1964)]; M. S. Ioffe, B. I. Kanaev, V. P. Pastukhov, V. V. Piterskii, and E. E. Yushamov, *Ibid.* 13, 1210 (1987) [Sov. J. Plasma Phys. 13, 697 (1987)].
[71]M. Haines, Nucl. Fusion 17, 811 (1977).
[72]G. F. Abdrashitov, A. A. Berkhtenev, V. V. Kubarev, V. E. Pal'chikov, V. I. Volosov, and Yu. N. Yudin, see Ref. 63, page 335.
[73]V. M. Panasyuk and F. A. Tsel'nik, Fiz. Plasmy 1, 552 (1975) [*sic*].
[74]B. Lehnert, Phys. Scripta 13, 317 (1976).
[75]V. N. Bocharov, N. A. Zavadskiĭ, A. V. Kiselev, S. G. Konstantinov, A. V. Kudryavtsev, O. K. Myskin, V. M. Panasyuk, and F. A. Tel'kin, Pis'ma Zh. Eksp. Teor. Fiz. 41, 494 (1985) [JETP Lett. 41, 601 (1985)].
[76]A. I. Morosov, see Ref. 4, Vol. 2, p. 3.
[77]V. V. Mirnov and D. D. Ryutov, Preprint IYaF SO AN SSSR 69–71 Novosibirsk, 1971; Nucl. Fusion 12, 627 (1972).
[78]R. F. Post, Phys. Rev. Lett. 18, 232 (1967).
[79]J. L. Tuck, see Ref.4, Vol. 2, page 595.
[80]G. I. Budker, V. V. Danilov, E. P. Kruglyakov, D. D. Ryutov, and E. V. Shun'ko, Pis'ma Zh. Tekh. Fiz. 6, 117 (1980) [Sov. Tech. Phys. Lett. 6, 52 (1980)]; Zh. Eksp. Teor. Fiz. 65, 562 (1973) [Sov. Phys. JETP 38, 276 (1973)].
[81]E. P. Kruglyakov, *Abstract of Doctoral Dissertation* [in Russian], IYaF SO AN SSSR, Novosibirsk, 1975.
[82]B. G. Logan, I. G. Brown, A. J. Lichtenberg, and M. A. Lieberman, Phys. Rev. Lett. 29, 1435 (1972); Phys. Fluids 17, 1302 (1974).
[83]H. D. Price, A. J. Lichtenberg, M. A. Lieberman, and M. Tuszewski, Nucl. Fusion 23, 1043 (1983).
[84]F. Najambadi, A. J. Lichtenberg, and M. A. Lieberman, Nucl. Fusion 23, 609 (1983).
[85]B. A. Knyazev and P. Z. Chebotaev, *Ibid.* 24, 555 (1984).
[86]A. V. Arzhannikov, B. N. Breizman, A. V. Burdakov, V. S. Burmasov *et al.*, see Ref. 60, p. 347.
[87]B. A. Knyazev, V. V. Mirnov, and P. Z. Chebotaev, *Problems in Atomic Science and Technology. Ser. Thermonuclear Fusion* [in Russian], TsNIIatominform, Moscow 3, 12 (1983).
[88]V. V. Mirnov and D. D. Ryutov, Pis'ma Zh. Tekh. Fiz. 5, 678 (1979) [Sov. Tech. Phys. Lett. 5, 279 (1979)].
[89]P. A. Bagryanskij, A. A. Ivanov, V. V. Klesov, Yu. L. Koz'minykh, A. Kotel'nikov, Yu. I. Krasnikov, A. A. Podyminogin, A. I. Rogozin, G. V. Poslyakov, and D. D. Ryutov, see Ref. 58 3, 467.
[90]K. L. Lam, B. J. Leikind, A. Y. Wong, G. Dimonte, A. Kuthi, I. Olson, and H. Zwi, Phys. Fluids 29, 3433 (1986).
[91]V. V. Mirnov and D. D. Ryutov, see Ref. 87 1, 57 (1980).
[92]A. I. Kotel'nikov, V. V. Mirnov, V. P. Nagornyj, and D. D. Ryutov, see Ref. 60, page 309.

[93]*International Tokamak Reactor*: Executive Summary by INTOR Group, Nucl. Fusion **25**, 1791 (1985).
[94]TASKA, A Tandem Mirror Fusion Engineering Test Facility: KFK-reports 3311/2 and UNWFDM-500 **1**, 2 (1982).
[95]*A Tandem Mirror Fusion Engineering Test Facility*, Preprint UCID-19328, Lawrence Livermore Nat. Lab., 1983.
[96]V. V. Mirnov, V. P. Nagornyĭ, and D. D. Ryutov, Preprint IYaF SO AN SSSR, No. 84-40, Novosibirsk, 1984.
[97]T. Kawabe, S. Hiroyama, Y. Kozaki, K. Yoshikawa *et al.*, Fusion Techn. **2**, 1 (1986).

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