The ‘SXR—Pit’ paradox as an indicator of electron transport in tokamaks

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Abstract
A phenomenon is analyzed, which is observed at some laboratory-scale tokamaks and seems to be unexplained: a strong drop in the soft x-ray radiation (SXR) after it passed through the ‘thick’ Be foils (the pit in the SXR spectrum). This phenomenon becomes more pronounced as the foil thickness increases and the plasma electron density decreases. The explanation for this phenomenon is proposed, which involves the assumption that the Maxwellian electron velocity distribution becomes ‘depleted’ in the range of velocities that are 3–5 times higher than the thermal one. The anomalous electron heat transport along the weakly perturbed toroidal magnetic field of the tokamak (the ‘magnetic flutter’ model) seems to be the most probable reason for this depletion. Thus, the ‘SXR—Pit’ can be a new tool for studying the physical nature of the anomalous electron transport in tokamaks.

Keywords: tokamak, lithium, SXR, electron confinement, magnetic flutter

(Some figures may appear in colour only in the online journal)

1. Introduction. ‘SXR—Pit’ phenomenon

In laboratory experiments with ohmic heating (OH) at some Russian tokamaks (the T-4, T-3M, and T-11M), the term ‘SXR—Pit’ denoted the phenomenon of abnormally strong absorption of the soft x-ray radiation (SXR) plasma radiation, measured using the foils method (e.g. [1]), under conditions of relatively low plasma electron temperature (∼200–400 eV) and low electron density $n_e$ (∼$10^{19}$ m$^{-3}$ or even lower). The foils method is known to be based on the effect of absorption of SXR radiation (1–10 keV) by several foils of different thicknesses (most often, a set of Be foils is used). The anomalous drop in absorption is especially pronounced when using ‘thick’ foils (e.g. 90 µm Be). Without going into detail, we note that this incomprehensible phenomenon was considered to be a manifestation of some complex behavior of impurities under conditions of the anomalous electron heat conduction at low densities $n_e$. However, in experiments with the lithium limiter at the T-11M tokamak, performed under conditions of pure plasma ($Z_{\text{eff}}(0) \sim 1$), the SXR—Pit effect became highly pronounced (see, e.g. [2]); although, in these experiments, the effect of heavy impurities was almost excluded. That is, when trying to explain the SXR—Pit phenomenon, it was no longer possible to refer to the impurities’ influences. (Some distinctive features of the ‘lithium experiment’, which made possible the T-11M operation at plasma densities lower than $10^{19}$ m$^{-3}$ without evidence of the runaway electron beam development are described in the appendix).

It is known that in pure hydrogen plasma with a fixed electron temperature, the SXR emission of electrons should be determined by their bremsstrahlung and, therefore, it should depend on the density as $n_e^2$. Therefore, with decreasing density $n_e$, when such a decrease in the SXR emission intensity is observed, it is often taken for granted and the fact is ignored that, in tokamaks in the OH regimes, a decrease in the electron density is usually accompanied by an increase in the electron temperature. The latter is clearly indicated by, e.g. an increase in the Spitzer electrical conductivity of...
the plasma column and a decrease in the loop voltage $V_L$, which is the parameter inverse to the electrical conductivity. In the absence of impurities and runaway electrons, the Spitzer electrical conductivity of the plasma is determined by the electrons of the Maxwell distribution in the thermal energy region ($\sim T_e$) and strictly follows $\sim T_e^{3/2}$; that is, it can be used as a primary indicator for estimating the energy distribution of electrons near $T_e$. Figure 1 shows an example of a change in the SXR signal that passed the 30 $\mu$m Be filter as the average $N_e$ density decreased on a T-11M tokamak [2] with a lithium limiter, when, as mentioned above, the influence of impurities and runaway electrons (appendix) could be ignored.

From figure 1 it follows that the intensity of SXR decreases approximately as $N_e^{-2}$. However, the SXR emission intensity strongly depends on the electron temperature (as the temperature increases, e.g. from 300 to 400 eV, the SXR emission intensity, after passing through the 90 $\mu$m Be foil, should increase approximately three times). With allowance for this fact, we would expect that the detected SXR radiation will even increase with decreasing $N_e$, or its decrease (proportional to $N_e^{-2}$) will become considerably slower. However, in experiments at the T-11M, no decrease was observed in the rate of SXR intensity drop; moreover, the SXR signal decreased even slightly faster compared to the function $\sim N_e^{-2}$. To eliminate this contradiction, the author analyzed in more detail time evolution of the SXR signal at the T-11M in the range of low plasma densities. In this analysis, the same database was used as in [2], namely, the T-11M tokamak shots #13 100–13 133.

In addition, the data on the SXR measurements performed using the 30 $\mu$m Be foils were supplemented by those using the 90 $\mu$m Be foils and earlier results of laser measurements of $T_e$ in similar modes of discharge of the tokamak T-11 [3], confirming the growth of $T_e$ from 300 eV to 400 eV during the density decrease from $5 \times 10^{19}$ m$^{-3}$ up to $0.7–1 \times 10^{19}$ m$^{-3}$. Next, to weaken the effect of the strong density dependence of the SXR signal ($\sim N_e^{2.5}$) and, thereby, to make the comparison easier, the results of both measurements were divided by $N_e^{2}$. In addition, the results of measurements with the 90 $\mu$m Be foils were multiplied by 5 or 4, which, at an electron temperature of approximately 300 eV (5) and 400 eV (4), typical of the T-11M tokamak, approximately compensates for the calculated excess of the signals transmitted through the 30 $\mu$m Be foils over those transmitted through the 90 $\mu$m Be foils. The comparison of the results finally obtained is shown in figure 2, in which the electron density $N_e$ (averaged over the central chord) varies approximately ten times (from 0.3 to $4 \times 10^{19}$ m$^{-3}$). Squares correspond to the passage of SXR 30 $\mu$m Be and triangles to the 90 $\mu$m Be foils. The curve A–A denotes the expected growth of the SXR radiation that has passed through the 30 $\mu$m Be foils as the $T_e$ increases from 300 eV to 400 eV, obtained on the basis of laser scattering measurements of $T_e$. Note that in the range of $N_e$ from 1.4 to $4 \times 10^{19}$ m$^{-3}$, squares and rectangles almost merge, which means that the $T_e$ values measured by laser scattering and the foils method are practically equal. That is, in this range of $N_e$ changes, the distribution of plasma electrons by energy should correspond to the Maxwellian one (this is usually assumed in the foils method).

However, below $1.5 \times 10^{19}$ m$^{-3}$, the course of the SXR signals is noticeably separated from the AA curve. If the deviation of the SXR signals that passed through the 30 $\mu$m Be foil could still be explained by experimental errors, then at $N_e < 1 \times 10^{19}$ m$^{-3}$, the signal of SXR that passed through the thick foil drops almost twofold, which formally could mean a decrease in $T_e$ from 400 eV to almost 250 eV. At the same time, the average $\mu$ in the absence of impurities and runaway electrons ($N_e > 1 \times 10^{19}$ m$^{-3}$) could still be explained by experimental errors, then at $N_e < 1 \times 10^{19}$ m$^{-3}$, the signal of SXR that passed through the thick foil drops almost twofold, which formally could mean a decrease in $T_e$ from 400 eV to almost 250 eV. At the same time, the average $\mu$ in the absence of impurities and runaway electrons (appendix) could be ignored.

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time, however, the Spitzer electrical conductivity of the plasma did not fall, but even slightly increased as \( N_e \) decreased.

This seemingly dead-end situation inherently discredited the idea of measuring the electron temperature in tokamaks at low plasma densities using the foils method. In laboratory slang, this situation was called the SXR—Pit occurrence, as a kind of warning to those who were planning to measure the plasma temperature using the foils method at \( n_e < 1 \times 10^{19} \text{ m}^{-3} \).

The threshold conditions for the transition to the regime with the SXR—Pit depend not only on the density, but also on the discharge current. All the experiments described above (shots ## 13 100–13 133) were performed in a relatively narrow range of operating parameters of the T-11M tokamak: the discharge current was approximately 75–85 kA, and the magnetic field was 1–1.2 T with the stability margin \( q(a) \sim 3.5–4 \). It turned out that, by reducing the discharge current to 60 kA and thereby increasing the stability margin up to 5, it was possible to lower the threshold density \( N_e \) for the occurrence of the SXR—Pit and, in contrast, by increasing the current to 90 kA, the threshold density can be made higher.

It is stated below that all these seemingly paradoxical properties of the SXR—Pit formation could be explained, if we assume the ‘depletion’ of the far ‘tails’ of the Maxwellian electron velocity distribution, as a result of the action of some specific mechanism for the high-energy electron loss along the main magnetic field of the tokamak.

Let us make estimates using the data of the T-11M experiment with the lithium limiter and the almost pure hydrogen plasma in the center of the plasma column, in which, as was stated above, the effect of impurities could be neglected. By analyzing the transmission of the SXR radiation through 90 \( \mu \text{m} \) and 30 \( \mu \text{m} \) Be foils, it was possible to estimate the energy range in the ‘tail’ of the Maxwellian electron energy distribution, corresponding to electrons, which, in the absence of depletion, could provide the SX radiation in the required energy range. Figure 3, assuming a Maxwellian electron distribution, shows the calculated convolutions of three functions: bremsstrahlung at \( T_e = 400 \text{ eV} \), and the carrying capacity of Be foil with a thickness of 30 and 90 \( \mu \text{m} \) for the energy range of interest from 1.4 to 6 keV [1].

As follows from the above figure, the SXR transmitted through the 90 \( \mu \text{m} \) Be foils for the plasma temperature of 400 eV of interest should be generated by electrons mainly in the range of 2500–1500 eV. On the other hand, SXR transmitted through the 30 \( \mu \text{m} \) Be foils should be generated by electrons mainly in the range of 2000–1200 eV. If we assume that the cause of the ‘anomalous’ attenuation of the SXR signal that passed through the 90 \( \mu \text{m} \) Be foils is the depletion of the Maxwell electron distribution, it is logical to expect that it will manifest itself mostly noticeably in the energy range of 2500–2000 eV with a possible deepening as it approaches 2500 eV. The depth of depletion of the Maxwellian distribution in this energy region should then be determined by the balance of its ‘filling’ due to collisions and the ‘drain’, which is an additional channel for the loss of energetic electrons under the influence of some mechanism of their departure, specific to the tails of their Maxwellian distribution.

It is believed that the Maxwellian distribution of electrons is established as a result of the Coulomb electron–electron collisions within the characteristic times \( \tau_e = 2.3 \times 10^3 T_e^{3/2}/n_e \) [4] \((Z = 1)\). In the range of plasma parameters under consideration, the characteristic time \( \tau_e \) is approximately 10 \( \mu \text{s} \) and, therefore, it was considered as an axiom that the Maxwellian equilibrium of electrons is set in tokamaks.

The appearance of the SXR—Pit indicated the obvious violation of this axiom in the OH regimes of the T-11M tokamak with low plasma densities, which necessitates the involvement of an additional mechanism that provides anomalous electron losses in the energy range near 2500–2000 eV during the discharge. Could it not be connected to the mechanism determining the effect of the so-called ‘anomalous electronic thermal conductivity (heat transfer)’, which is most clearly manifested in the same region of small \( n_e \)?

The consequences of occurrence of the anomalous electron thermal conduction are well known (in terms of energy loss). Their quantitative measure is the magnitude and dynamics of the plasma energy confinement time \( \tau_E \), defined as a ratio of the total plasma energy to the heating power (as a rule, in the OH regime, the tokamaks operate in the quasi-stationary mode, when the powers of losses and heating are almost equal). Strictly speaking, the energy confinement time is characteristic of the energy losses of the entire plasma (both ions and electrons), but in the range of low densities \( n_e \), which is important to us, the contribution of ions is small and, with an accuracy of 20%–30%, the plasma energy confinement time in OH regimes is equal to the electron energy time, \( \tau_E \approx \tau_{Ee} \). Accepting \( \tau_E \) as a quantitative characteristic of the anomalous electron thermal conduction and assuming that both the anomalous electron transport and the escape of high-energy
electrons from the plasma column occur under the influence of the same mechanism, we can estimate the limiting rate of depletion of the Maxwellian electron energy distribution in the energy range corresponding to the tails, and compare it with the rate of ‘establishment’ of the Maxwellian electron energy distribution in the same energy range, as a result of the Coulomb electron–electron collisions.

2. The Coulomb ‘establishment’ of the Maxwellian electron energy distribution

It is generally accepted that the Maxwellian distribution at energies $E_e \sim T_e$ is already established after one or two electron–electron collisions. To estimate the time required for its establishment at the level $E_e \sim 10 T_e$, we will use the results of the analytical and numerical calculations performed in [5] by Potapenko and co-authors. Figure 4, as development of [5], shows the correction factor to the Maxwellian distribution, $g(v; \tau) = f(v; \tau) f_{ee}(v)$, calculated in accordance with the calculations of the Coulomb interaction potential, as a function of the relative velocity of particles participating in elastic collisions (in units of thermal velocity $v_{th}$, at $E_e \sim T_e$; more precisely, $m_e v_{th}^2 = 3 T_e$) and the number of collisions (in arbitrary units) also at $E_e \sim T_e$. From the given figure, it follows, in particular, that already two collisions are sufficient for the establishment of the Maxwellian distribution (the correction factor is $g(v; \tau) = 0.8$) for electrons with energies less than $E_e = 4 T_e$. If we consider the level of practical filling of the Maxwell distribution $g(v; \tau) = 0.8$, then eight collisions (figure 4) would be enough to fill it up to 3200 eV—near the upper limit of the energy range of interest to us. If we take into account that the main part of the energy range that is important to us is located somewhat lower, in the zone of 2500–2000 eV, the number of necessary collisions should be reduced to 6–5.5, in proportion to $(v_{min}/v_{th})^3$ [5]. We will take this level as a fill indicator, below which the depletion of the Maxwell distribution in the energy range of interest will become noticeable, and the corresponding time $(6/v_{ee})$ will be considered a characteristic time for establishing the Maxwell distribution in this energy range.

How could anomalous electronic heat transfer provide for the depletion of the energy interval $E_e \approx 2500–2000$ eV at $T_e = 400$ eV and $N_e \sim 10^{19} \text{ m}^{-3}$ in a time of $6/v_{ee}$ or shorter?

3. Anomalous electron thermal conduction (heat transport)

The term ‘anomalous electron heat transport’ indicating the main channel for the tokamak energy losses, without specifying its physical nature, in contrast with official ‘anomalous ion heat transport’ was first introduced by Artsimovich in 1968 in the review report at the 3rd IAEA International Conference on Plasma Physics and Controlled Nuclear Fusion Research (Novosibirsk, 1968) [6] after it was experimentally discovered at the T-3 tokamak that the plasma energy confinement time $\tau_E$ increases with increasing density (approximately as $n_e^{0.3}$). In this case, in the low density range ($<1 \times 10^{19} \text{ m}^{-3}$), the lifetime $\tau_E$ turned out to be much less than the calculated time of the electron–ion energy exchange, that is, by no means, it could be determined by the ion energy losses. This was the drastic violation of the basic concept of the magnetic thermal insulation of plasma in tokamaks, within which, as is known, the energy confinement time of hot plasma should be determined by the ion losses and decrease with increasing density (and, accordingly, the collision rate). But it turned out that, in the OH tokamaks, the measured energy confinement time of electrons varies the other way around: it increases with increasing plasma density and collision rate, just as would be expected, for example, in the inertial fusion facilities (the H-bomb model). Figure 5 [6] shows the shocking dependence of the plasma energy lifetime on the density measured at the T-3 tokamak, which turned out to be a ‘pain spot’ for physicists of that time, as well as for a number of physicists of later times.

As can be seen in figure 5, at the T-3 tokamak, the steepest drop in the energy confinement time $\tau_E$ (almost linear with decreasing $n_e$) was observed in the same density range, in which, much later, the SXR—Pit effect was discovered at the T-11M tokamak (figure 2).

In supporters of the ‘classics’, for some time there were hopes that the basis of the measurement of the energy of the plasma on the tokamak T-3 (it was measured by the plasma diamagnetism [7, 8]) was made of some underlying methodological error. But in the heat of controversy that arose at the 3rd IAEA conference, Artsimovich suggested to British physicists that they come to Moscow to Tokamak-3 with their new laser technique of Thomson scattering and measure the main critical plasma parameter—the electronic temperature—which, following Russian diamagnetic measurements, was approaching 1 kV. They unexpectedly agreed to come to Moscow the following year, 1969, and were convinced of his correctness. The laser scattering [9] showed that the results of diamagnetic measurements on T-3 were deep within the given errors. Finally, after ten years, a linear growth of $\tau_E$ with a density in the widest range from 0.5 to $6 \times 10^{20} \text{ m}^{-3}$ was
Peripheral magnetohydrodynamics (MHD) turbulence, which, the authors of the idea, their primary cause could be, e.g. the limiter, or the divertor), under the effect of small disturbances during one collision. The particle can be displaced by \(2^\frac{\lambda}{\nu}\), the mean free path is \(\lambda = \nu\sqrt{m}\), where \(m\) is the particle mass, and \(\nu\) is its transverse velocity with respect to the magnetic field \(B\). As for the trap with the longitudinal plasma confinement (along the magnetic field \(B\)) (figure 6(B)), the mean free path is \(\lambda = \nu/v\), where \(\nu\) is the collision rate of identical particles.

Within the framework of the classical model of the pair of Coulomb collisions, the transport (diffusion and heat conduction) of high-temperature plasma should follow the well-known Einstein law:

\[
(\Delta x)^2 \sim t, \text{or } (\Delta x)^2 = Dt
\]  

where \(t\) is the spreading time of the particle ensemble, \(\Delta x\) is the linear scale of spreading, and \(D\) is the effective diffusion coefficient, or, in the case of heat transport, \(\chi\) is the thermal diffusivity.

We denote again that \(\lambda\) is the mean free path of particles, and \(\nu\) is the collision rate. With allowance for the facts that, during one collision \(t = 1/\nu\), the particle can be displaced by not more than \(\Delta x = \lambda\), and the order-of-magnitude estimate

4. Transverse and longitudinal confinement of plasma

To make the further presentation more illustrative, the schematics of the plasma facilities are shown in figure 6: (figure 6(A)) the classical tokamak, which is the facility with the transverse plasma confinement; and (figure 6(B)) the magnetic tube filled with plasma, resting upon the zero-temperature end (e.g. the limiter). Arrows indicate the directions of heat flows. The inscriptions on the arrows show their possible values calculated within the framework of the model of the classical plasma transport along and across the magnetic field (calculated according to Braginsky [4]). All the subsequent quantitative estimates are made by the author using the CGS system of units (to avoid discrepancy), as this system of units was used by Braginsky, the authors of [5], and in the well-known reference book [15]. In the CGS system of units, the temperature and energy are measured in eV, the length and velocity are measured in cm and cm s\(^{-1}\), respectively, and the density \(n_e\) is in the units of \(\text{cm}^{-3}\).

We briefly remind the reader what should be the most striking distinctive features of plasma in the trap with the transverse magnetic confinement (figure 6(A)). In such traps, the mean free path \(\lambda\) of charged particles is determined by the Larmor radius \(\rho = \lambda \sim mvB\), where \(m\) is the particle mass, and \(v\) is its transverse velocity with respect to the magnetic field \(B\). As for the trap with the longitudinal plasma confinement (along the magnetic field \(B\)) (figure 6(B)), the mean free path is \(\lambda = v/\nu\), where \(\nu\) is the collision rate of identical particles.

Among these attempts, the most consistent (in the author’s opinion) is the idea of the ‘longitudinal transport’ of electron heat as a result of the partial destruction of the tokamak magnetic configuration, e.g. the formation of the small magnetic islands, comparable in size with the electron Larmor radius, but smaller than the ion Larmor radius (Callen called them the ‘magnetic flutter’ [12], and Kadomtsev and Pogutse introduced the concept of the ‘braided magnetic field’ [13]). Such low-scale ‘damages’ to the tokamak magnetic configuration \(\delta\) cannot radically change the ion transport across the magnetic field. But, at the same time, their effect is quite sufficient to transform the classical transverse electron transport onto the wall into the longitudinal one: the transport along some assumed magnetic field line with the length \(L\), which winds many times around the torus and, with each revolution, gradually approaches the chamber wall (or the limiter, or the divertor), under the effect of small disturbances \((b_\tau)\) of the transverse magnetic field. As schemed by the authors of the idea, their primary cause could be, e.g. the peripheral Magnetohydrodynamics (MHD) turbulence, which, in the central regions of the plasma column, induces microperturbations of the magnetic configuration with the transverse size not exceeding the size of the collisionless skin layer \(c/\omega_{pi}\) (the ‘Ohmawa parameter’ [14]). Thus, for electrons, which, in the OH regimes with low densities \(n_e\) in tokamaks, are the main carriers of the thermal plasma energy, the energy transport would become the longitudinal one. What fundamental differences in the plasma macroscopic behavior would occur in this case, and why does the idea of the longitudinal transport look the most convincing for the author, despite its seeming paradoxality?

**Figure 5.** The T-3 tokamak [6]. Dependence \(\tau_E(n_e), J_p = 40\,\text{kA},\) and \(B_T = 2.5\,\text{T}\).
of the $D$ coefficient is $\lambda^2 \nu$, we can rewrite dependence (1) as follows:

$$ (\Delta x)^2 = \lambda^2 \nu t. $$

After we assume $\Delta x = a$, or $\Delta x = L$, this universal equation can be used as the order-of-magnitude estimate of the energy confinement time $\tau_E = t$ in the magnetic traps with both the transverse confinement (figure 6(A)) and the longitudinal confinement in the magnetic tube with the length $L$ (figure 6(B)).

Within the framework of the model of the pair of Coulomb collisions, we make the order-of-magnitude estimate of the characteristic times for the hot plasma confinement in such facilities. We note that in this case, the collisions of particles of the same type ($e-e$, or $i-i$) will determine the corresponding heat conduction, while the collision of particles of different types ($e-i$) will determine diffusion. For definiteness, we restrict ourselves to the heat conduction estimates, since this problem is the closest to the problem under consideration.

In the trap shown in figure 6(A), the heat flow from the plasma falls onto the chamber wall with radius $a$ and, in the cylindrical geometry, the effective free path should be equal to the Larmor radius $\lambda_L$ (in the case of the closed magnetic traps, it is $q \lambda_L$). Whence it follows that the expected energy lifetime $\tau_E$ (the characteristic time, during which the heat reaches the wall) should be proportional to $a^2 (\lambda_L)^2 \nu_i$ and $a^2 (\lambda_L)^2 \nu_e$ for ions and electrons, respectively; here $\nu_i = 1/\tau_i$, and $\nu_e = 1/\tau_e$.

After performing the substitution, in the cylindrical geometry, we obtain the well-known classical expressions for $\tau_E$: for ions,

$$ \tau_{Ei} = \frac{a^2 T_i^{1/2} B^2}{n_i m_i^{1/2}}. $$

and for electrons:

$$ \tau_{Ee} = \frac{a^2 T_e^{1/2} B^2}{n_e m_e^{1/2}}. $$

Hence it follows that, for approximately equal temperatures, $T_i$ and $T_e$, and also densities, $n_i$ and $n_e$, $\tau_{Ee}$ should approximately 40 times exceed $\tau_{Ei}$. In tokamaks, taking into account the trapped particles [16] would reduce the indicated lifetimes $\tau_E$ by approximately $(B_p/B)^2$ times (where $B_p$ is the magnetic field of the current $J_p$), but the functional dependencies and, accordingly, the correlations between $\tau_{Ei}$ and $\tau_{Ee}$ will remain unchanged.

In the trap shown in figure 6(B), the heat flow from plasma is limited in the transverse direction by the magnetic field $B$, so the heat flows along the magnetic field to the end $S$. In this case, the length $L$ of the corresponding ‘plasma tube’ becomes the linear limit for the ‘ensemble spreading’, and the mean free paths for ions and electrons are $\lambda_i = v_i \tau_i$ and $\lambda_e = v_e \tau_e$, respectively; they are the same at equal temperatures and densities. After performing the substitution, we obtain the well-known classical expressions for $\tau_E$ in the case of the longitudinal transport, namely:

for ions:

$$ \tau_{Ei} = \frac{L^2 n_i m_i^{1/2}}{T_i^{3/2}}. $$

and for electrons:

$$ \tau_{Ee} = \frac{L^2 n_e m_e^{1/2}}{T_e^{3/2}}. $$

Hence, it follows that, under conditions of the longitudinal transport, the relation between the energy lifetimes becomes
inverse compared to the case of the transverse transport: the electron transport should approximately 40 times exceed the ion transport. In addition, if we take into account that in OH conditions, $T_i < T_e$, the ratio of energy lifetimes can become 100 or even higher.

We note that, at low plasma densities, such a longitudinal energy transport would mostly affect the high-energy electrons, which qualitatively coincides with the observed formation of the SXR—Pit. To what extent could it be responsible for its occurrence? Let us make the quantitative estimates. We compare the aforementioned Coulomb mechanism for calculating the ‘filing’ of the Maxwellian electron energy distribution with the assumed mechanism for the longitudinal depletion. The comparison will be performed for electrons with energies of $\sim 2.5$ keV under real conditions of the T-11M tokamak, assuming the longitudinal electron heat conduction to be the dominant cooling channel for the plasma electron component (in accordance with calculations of Braginsky [6]) (figure 6(B)). First, we consider the plasma in the standard regime of the T-11M with temperature $T_e$ and density $n_e$, for which the energy lifetime is known, $\tau_{Ee} \approx T_e$, that is, as suggested above, we neglect the contribution of the ion component, as it is relatively small. Let us make the order-of-magnitude estimate of energy flows from the magnetic tube with the length $L$ and the end (limiter) $S$, filled with plasma with the density $n_e$ and the average temperature $<T_e> \sim T_e/2$, shown in figure 6(B). If we assume that $\nabla \cdot T_e \approx T_e/L$, then, under conditions of equilibrium, the electron heat flow, coming into the plasma from the left and then falling onto the tube end, $S^e_\parallel T_e \approx S^e_\parallel T_e/L$, should compensate for its energy losses, namely, $W_\parallel/\tau_{Ee}$. Here, $W_\parallel$ is the electron thermal energy of the plasma filling the tube, which, under the assumption that the number of electron–electron (e–e) collisions is large enough ($L > \lambda_{ee}$), is equal to $0.75n_eL_eSL$. The electron energy lifetime $\tau_{Ee}$ of such plasma can be written in the following form:

$$\tau_{Ee} = 0.75n_eL_e^2/T_e \chi_\parallel$$

(7)

or, with allowance for expressions (8) and (9),

$$\chi_\parallel = 3.16n_eT_e\tau_e/m_e$$

(8)

(see [3]),

and $\tau_e = 2.3 \times 10^4 T_e^{3/2}/n_e$ (at $Z = 1$),

(9)

we obtain the following expression:

$$\chi_\parallel = 3.16n_e T_e \tau_e/m_e \approx 1.5 n_e \tau_e v_{th}^2$$

(10)

where $v_{th} = (2kT_e/m_e)^{1/2}$ is the average thermal velocity of the electron component.

Next, substituting expression (7) into expression (10), we obtain the following simple estimate for the electron energy time $\tau_{Ee}$:

$$\tau_{Ee} \approx L^2 \left( \tau_e v_{th}^2 \right)^{-1}.$$  

(11)

If we take into account that the electron energy confinement time $\tau_{Ee}$ obtained in this way should be approximately equal to the total plasma energy time $\tau_p$ measured experimentally by dividing the total plasma energy by the heating power, then it will be possible to estimate the length $L$ of the assumed magnetic field line, connecting the hot region of the plasma column and the first wall of the tokamak (or rather limiter), as follows:

$$L \approx v_{th} (\tau_p \tau_e)^{1/2}. $$

(12)

How long can $L$ be in T-11M? We analyze two T-11M shots shown in figure 2, namely, shot #13130 ($N_e \approx 4.5 \times 10^{15}$ cm$^{-3}$ and $\tau_E \approx 5 \times 10^{-3}$ s at $T_E(0) \approx 270$ eV), which is the high-density shot, shown on the right-hand side of the figure and corresponds to the ‘normal’ density range far from the SXR—Pit, and shot #13117 ($N_e \approx 1 \times 10^{13}$ cm$^{-3}$ and $\tau_E \approx 2 \times 10^{-3}$ s at $T_E(0) \approx 400$ eV), which corresponds to the SXR—Pit formation.

Assuming again $\tau_e = 2.3 \times 10^4 T_e^{3/2}/n_e$ (\Lambda = 15), we obtain for shot #13130 in the normal density range: $L = 1.2 \times 10^3 m$ at $\lambda_{ee} = 30$ m, that is, the number of collisions $k$ of the ‘thermal’ electron during its supposed drift to the limiter turns out to be of the order of 40. That means that the thermal electron drift regime is deliberately collisional, and the Maxwellian electron energy distribution is undoubtedly established.

Next, we analyze the motion of the ‘fast’ electron component with energies $E_e = 2500$ eV along the same magnetic field line $L$ with allowance for the damping of the fast electrons in collisions with the electrons of the surrounding ‘warm’ (400 eV) plasma. We will use the expression for the characteristic time of collisions between the fast-electron energies with energy $E_e$ and the bulk of plasma with $T_e$, recommended in the Naval Research Laboratory (NRL) Plasma Formulary [15]:

$$\tau_e = 2.3 \times 10^4 T_e^{1/2} E_e/n_e (\Lambda = 15).$$

(13)

The characteristic mean free path of the fast electrons surrounded by the thermal electrons turns out to be approximately $5 \times 10^4$ cm, which is approximately 2.5 times shorter than $L$: that is, the energy exchange between the fast and thermal electrons should also occur in the almost collisional regime, which is ‘sufficiently collisional’ to establish the Maxwellian electron energy distribution in this energy range.

Next, using the same scheme, we analyze shot #13117 ($N_e \approx 1 \times 10^{13}$ cm$^{-3}$ and $\tau_E \approx 2 \times 10^{-3}$ s at $T_E(0) \approx 400$ eV) close to SXR—Pit formation (see figure 2).

As well as for shot #13130, we first consider the bulk of plasma: that is, the energy range $E_e \sim T_E(0)$. Assuming again that $\tau_e = 2.3 \times 10^4 T_e^{3/2}/n_e$ (9), we obtain that, for $N_e \approx 1 \times 10^{13}$ cm$^{-3}$ and $\tau_E \approx 2 \times 10^{-3}$ s, the length of the magnetic field line is $L = 1.6 \times 10^3$ m. In this case, it turns out that the mean free path is $\lambda_{ee} \approx 150$ m and, accordingly, the number of collisions $k$ is of the order of 10: that is, the further electron energy exchange in the range of thermal energies will be ensured by the collisions, and the Maxwellian electron energy distribution will be naturally established.
But in the electron energy range in the vicinity of 2500 eV, the situation is more complicated.

The characteristic free path length for these electrons should be $\approx 2.4 \times 10^5$ cm, which significantly exceeds $L$, and thus requires a more detailed analysis. The time flight of such electrons with energy $E_e = 2500$ eV from the hot plasma region to the limiter (the ‘escape’ of electrons) is too short. It will be equivalent to only about four collisions in warm plasma (figure 4), while 5–6 such collisions would be necessary to ‘restore’ the Maxwell distribution to the level of $g(v, t) \approx 0.9–0.7$ near its $\approx 2500$ eV. Thus, at the point—$N_e \approx 1 \times 10^{13}$ cm$^{-3}$ (shot # 13 117), we could find ourselves in the region of noticeable depletion of the Maxwell distribution in the region of electronic energies near $E = 2500$ eV. Thus, the formation of the SXR—Pit at the T-11M at low plasma densities could be naturally described within the framework of the model of longitudinal confinement of the plasma electron component [12, 13].

5. Discussion

5.1. SXR—Pit and magnetic flutter

To what extent do the conditions of the described experiments fit into the framework of the conditions imposed by the longitudinal confinement model of magnetic flutter [12, 13]? Namely, how does the transverse size of the expected ‘disturbances’ of the T-11M magnetic configuration (destruction length) correlate with the electron and ion Larmor radii?

Let us make the order-of-magnitude estimate of the expected destruction length $\delta_M$ of the magnetic configuration for shot #13 117 with $L = 1.6 \times 10^5$ m. Taking into account that, in the T-11M, the circumference of the torus is $2\pi R = 4.4$ m, we obtain that, during the transverse drift of the magnetic field line with the length $L = 1.6 \times 10^5$ m = $2\pi R N$ from the axis of the plasma column to the limiter, it should have made $N = 360$ turns along the torus. Taking into account that, at the T-11M, the minor plasma radius (the distance from the axis to the limiter) is approximately 17 cm, the transverse outward displacement $\delta_M$ of the field line along the radius $r$ per revolution around the torus (which can be accepted as a linear measure of the ‘perturbation’ of the tokamak magnetic configuration) should be approximately 0.5 mm. In this case, for $T_e = 350$–400 eV and $B_T = 1$ T, the electron Larmor radius turns out to be approximately ten times less than the displacement $\delta_M$ (0.05 mm), and the ion Larmor radius is four times larger (2 mm). That is, the main critical parameters wittingly comply with the requirements for the longitudinal confinement of electrons in the magnetic flutter model [8, 9].

What should be the radial component of the magnetic field perturbation $(\delta_B)$, averaged over the circumference of the torus, to provide such a radial displacement of the magnetic field line? It can be estimated that in this case, it will be sufficient, if the perturbation amplitude is approximately 0.1% of the magnetic field $B_p$ (the magnetic field of the current $I_p$). We remind the reader of the fact that the amplitudes of the magnetic disturbances accompanying, for example, the development of the tearing instability in tokamaks, are an order of magnitude or even higher: that is, we are considering the phenomenon of the next order of smallness.

An order of magnitude difference in the amplitudes of the magnetic disturbances indicates that at the T-11M, the energy store of the mechanism capable of transforming the transverse electron transport into the longitudinal one should be two orders of magnitude less than that maintaining the large-scale MHD activity in the plasma column.

At last, we compare the $\delta_M$ displacement with the parameter $c\omega_{pl}$ (the width of the collisionless skin layer), which characterizes the ‘permissible’ displacements introduced by the so-called Ohkawa formula [13, 14] for estimating the turbulent heat transport $(\chi \approx c^2 vre/\omega_{pl}^2 qR)$. It is assumed [12] that within these limits, the magnetic reconnection and stochasticization of the magnetic disturbances are permissible (the generation of the symmetric disturbance $b_t$ is prohibited by the condition $\text{Div}(b) = 0$). However, if we take into account that all these processes occur under conditions of the dominating magnetic asymmetry with respect to the major radius $R$, then it will no longer seem surprising that the spatially averaged asymmetric disturbance arises against their background, $(b_t) \approx 10^{-3} R_T$. In the density ranges $N \approx 1 \times 10^{13}$ cm$^{-3}$ and $4 \times 10^{13}$ cm$^{-3}$ (high-density range), the corresponding parameters are equal to $c\omega_{pl} \approx 1.7$ and 0.5 mm, respectively. That is, in the entire density range of the T-11M tokamak operation, the linear ‘disturbance’ of the order of $\delta_M = 0.5$ mm would be in the range of displacements ‘permissible’ for the development of the corresponding magnetic disturbances $(b_t)$.

5.2. Competition of transverse and longitudinal transport

In what other macroscopic manifestations of plasma dynamics in tokamaks does longitudinal transport (B) successfully compete with transverse transport (A)?

Let us again compare the (A) and (B) concepts of the plasma confinement (figure 6). What predictions do they make concerning the plasma behavior in tokamaks, and are these predictions true in reality?

(a) As already mentioned above, in case (A), the ion channel should be the dominant channel for the heat (and particle) transport onto the wall. The electron transport should be approximately 40 times weaker. In case (B), in contrast, the electron transport should dominate. In reality, the electron transport dominates in the low-density range and, accordingly, in the range of low collision frequencies (LOC) [6]. It becomes comparable to the ion transport only in the range of high collision frequencies (SOC). That is, the transfer of plasma thermal energy to the tokamak wall demonstrates a clear visible ‘dualism’.

(b) As a result of what was noted above, in case (A), the energy losses should increase with increasing plasma density, and the corresponding energy lifetime $\tau_E$ should decrease. In case (B), in contrast, the heat losses should
Cyclotron Resonance Heating (ECRH) pulse.

Figure 7. Time evolution of the average plasma density $n_e$ and the plasma electric potential at the radius $r \approx 2a/3$ at the T-10 tokamak [19]. The initial time corresponds to the beginning of the Electron Cyclotron Resonance Heating (ECRH) pulse.

decrease with increasing density. They decrease really and, accordingly, the total energy lifetime $\tau_E$ either increases with increasing density $n_e$, or reaches the plateau (figure 4).

(c) As the electron temperature increases, e.g. due to Electron Cyclotron Resonance (ECR) heating, the frequency of collisions should decrease, and in case (A), the energy lifetime should increase as $\tau_E \sim T_e^{-1/2}$. In case (B), it should decrease as $\tau_E \sim T_e^{-5/2}$. As shown experimentally, $\tau_E$ decreases in experiments on the axial ECR heating at the T-10 tokamak in the range of low and moderate $n_e$ [17]. At the same time, two ‘dead-end’ phenomena are observed that could not be explained within the framework of the (A) model: a decrease in the electron density $n_e$ in the center of the plasma column (‘density pump-out’ effect [18]) and the degradation of the negative electric potential $\varphi(0)$ at the center of the plasma column [19] up to positive. With the permission of the authors of [18], figure 7 taken from their work shows an example of degradation of the potential $\varphi(0)$ on the axis of the plasma column (the T-10 shot #62499). In this shot, the ECRH power was 1.1 MW and, during the shot, the electron temperature $T_e$ increased from 1 to 1. 8 keV. In this case, the potential was measured by analyzing the heavy ion beam passing through the plasma column [19]. Synchronously with the $\varphi(0)$ degradation, a considerable decrease in the electron density (density pump-out effect) was clearly observed, which, obviously, was followed by the escape of ions from the core region of the plasma column. Within the framework of the model (B), it is natural to associate the escape of ions with the action of the electric field arising as a result of the longitudinal escape of the heated electrons from the center. Under the action of this field, ions follow the electrons, thus, participating in the longitudinal outward transport. Most likely, the same phenomenon underlies the so-called ‘isotopic effect’: a decrease in the rate of the ion escape from the center observed with growing ion mass $m_i$.

(d) Now we consider the isotopic effect: the dependence of the plasma energy lifetime on the mass of the plasma-forming ion (H, D, or T). In case (A), it should be negative: that is, as the ratio $M$ of the ion mass to the proton mass increases, the ion escape from the plasma should become faster. After that, their thermal insulation should have deteriorated. In accordance with formula (3), the approximate dependence should be as follows: $M^{-0.5}$. In case (B), in contrast, the ion escape from the plasma should become slower, and the thermal insulation should be improved. The approximate dependence should be as follows: $M^{0.5}$ (see formula (5)). The thermal insulation improvement, which occurs with increasing ratio $M$, has been repeatedly observed at different tokamaks, but it remains either one of the ‘inexplicable miracles’ or ‘hard-to-explain anomalies.’ As is known, the improvement in thermal insulation was clearly demonstrated in experiments with the Deuterium Tritium (DT) plasma at the Tokamak Fusion Test Reactor (TFTR) tokamak. In these experiments, as the $M$ ratio increased from 1.9 to 2.7, the thermal insulation improved in accordance with the approximate law $M^{0.67}$ [20]. We note that the conditions of these TFTR experiments were as close as possible to the ‘fusion’ conditions: that is, the discharge current was up to 2.5 MA and the DT neutron yield was of the order of 10 MW.

(e) Finally, in case (A), an increase in the longitudinal magnetic field $B$ should radically improve the thermal insulation of plasma, namely, it should improve as $\sim B^2$ or, in the cylindrical geometry, as $\sim B_r^2$. However, the neoclassical effects in tokamaks lower the magnetic field $B$ from the $B_r$ level to almost the $B_p \sim aJ_p$ level; nevertheless, the quadratic dependence $B_p^2 \sim (aJ_p)^2$ should have been preserved. In case (B), the thermal insulation should not depend on the longitudinal magnetic field, as well as on the current $J_p$. It turns out that the real plasma behavior corresponds to a certain concept, which is intermediate between the (A) and (B) concepts. As it was first demonstrated at the T-3 tokamak [6, 12, 21], in the ‘stable regimes’ ($\varphi(a) > 3$), the energy confinement time $\tau_E$ increased almost proportionally to the current $J_p$ and did not depend on the field $B_r$. In this case, the stable regime was characterized by the suppression of large-scale MHD disturbances (the ‘Mironov oscillations’), which occurred as the magnetic shear of the tokamak increased. The latter occurred as a result of the inward diffusion of the current $J_p$ into the plasma region with the higher temperature and, accordingly, the higher electrical conduction [22]. In plasma of the TM-3 tokamak, which was smaller in size, similar distinctive features were observed, and the corresponding energy confinement times $\tau_E$ were approximately proportional to the squared transverse size of the plasma column.

5.3. ‘Dualism’ of energy scaling laws of ‘closed systems’ (tokamaks and stellarators)

At the same 3 IAEA conference in 1968, Artsimovich did not very confidently propose a general formula describing the maximum achieved energy confinement time $\tau_E$ [6] in T-3
and TM-3 tokamaks. Combining the T-3 and TM-3 results, he obtained one of the possible scaling laws for $\tau_E$ only ‘for the purpose of making interpolation estimates in the steady state operating tokamak regime’ (in the main Russian version of his communication). ‘Not very confidently’ means the new scaling law was not what was expected of him.

The similarity law was written in the following phenomenological form:

$$\tau_E \sim a^2 B_p n_e^{1/3}$$  \hfill (14)

which later turned out to be the first ‘working’ $\tau_E$-scaling law for OH tokamaks, which adequately describes the maximum plasma parameters obtained at the new facilities with different geometric and magnetic properties. This scaling law changed several names until it became clear that this simple ‘interpolation’ formula (14) predicted quite accurately the maximum achievable energy times $\tau_E$ for all round OH tokamaks put into operation after the T-3 tokamak from the T-4 and further to T-10 and Princeton Large Toroid (PLT).

Finally, Kadomtsev respectfully renamed this formula, calling it the ‘T-4 scaling law’, and wrote it in the following form [23]:

$$\tau_{\text{E}98} \sim 1.5 \times 10^{-8} B_p <n_e>^{0.5} a^2 s \quad (15)$$

where $B_p$ is in units of Oe, $<n_e>$ is in units of $10^{13}$ cm$^{-3}$, and $a$ is in cm.

We note that, in this case, the factor $a^2$, which takes into account the difference in sizes of the T-3 and TM-3 (they were geometrically similar, $R/a \approx 4-5$), was introduced by Artsimovich somewhat arbitrarily, by analogy with the transverse energy transport, which, as it turned out later, was not quite correct.

This incorrectness was first manifested in experiments at the T-11 tokamak with the geometrical parameter $R/a = 2.8$ [24]. It turned out that a decrease in the plasma radius $a$ had almost no effect on the lifetime $\tau_E$. Comparison of databases of the T-3 and other tokamaks allowed the authors of [24] to propose the refined scaling law $\tau_E \sim R^{2.75} a^{0.25} n_e$ (the Merezhkin-Mukhovatov scaling law). This scaling law radically contradicted the idea of the diffusion electron transport across the magnetic field. Next, so-called ‘neo-Alcator’ scaling law [25], $\tau_E \sim R^{2} a n_e$, was obtained, which described time evolution of the lifetime $\tau_E$ at the Alcator tokamak with the strong magnetic field and intensive injection of neutral hydrogen during the operating shot. This scaling law also contradicted the idea of the transverse confinement (maybe, the contradiction was less radical, but quite clear).

As is known, 20 years later, the results obtained at more than ten tokamaks with different sizes, geometries, and heating methods (and heating powers $P_H$) operating during this period were summarized by the ITER experts, and the more detailed and complex scaling law for the energy lifetime $\tau_E$ was proposed not only for the ohmic heated, but also for all other tokamaks with the additional plasma heating power $P_H$ (H-mode, [21]):

$$\tau_{\text{E}98} = 0.0365 J_p^{0.07} B^{0.08} P_H^{-0.63} n_e^{0.41} M^{0.20} R^{1.91} (a/R)^{0.23} k^{0.67} s \quad (16)$$

where $M$ is the isotopic factor, and $k = b/a$ is the elongation of the plasma cross section in the vertical direction.

It seems that scaling laws (14) and (15) are similar, but they could not be compared, since scaling law (15) includes the new indeterminate factor $P_H$ (the additional heating power). However, this difficulty can be overcome, if we take into account that all scaling laws are usually compiled from data of the ‘high performance’ operating regimes, with the maximum permissible power $P_{HM}$. As was noted in [26], in these regimes, the following rather rigid condition is satisfied: the maximum permissible heating power $P_{HM}$ remains approximately proportional to the plasma-facing surface area of the tokamak discharge chamber, $S \approx 4\pi R^2 a^{1/2}$. Thus, for the model group of the high-performance operating regimes (and we note again that, as a rule, the discharge regimes on which scaling law (15) was based were the ‘best’ ones), the uncertainty associated with the $P_H$ power can be considerably reduced. Namely, for this group of the high-performance operating regimes, it is possible to write the following approximate scaling law:

$$\tau_{\text{E}98} \sim J_p n_e^{0.4} R^{1.2} a^{-0.5} k^{0.4} B_p n_e^{0.4} \quad (17)$$

or, with allowance for the correlation $J_p \sim B_p a^2$:

$$\tau_{\text{E}98} \sim R^{1.2} a^{-0.5} k^{0.4} B_p n_e^{0.4} \quad (18)$$

which seems to differ from the scaling law $\tau_E \sim a^2 B_p n_e^{1/3} (14)$ only by the ‘geometric’ factor $R^{1.2} k^{0.5}$ instead of $a^2$.

It is natural to assume that the geometric factor should be more ‘conservative’ than others. Namely, since it is valid for the high-performance regimes, it is likely to be valid for a wider range of phenomena associated with the plasma transport in tokamaks. In essence, the parameter $R$ is a geometric measure of the longitudinal transport to the same extent, as the plasma radius $a$ is the measure of the transverse transport. The point is that the length $L$ of the assumed magnetic field line connecting the hot core of the plasma column and its edge, which is in contact with the wall, should be $\sim R(b_i)$, where, as defined above, ($b_i$) is the radial component of the magnetic perturbation averaged over the circumference of the torus, which induces the longitudinal transport of electron energy in a trap of the B type. If it is approximately constant in the different operating regimes of tokamaks, the energy lifetime $\tau_E$ will vary as $\sim R^2$ (in case B of the purely longitudinal transport) or $\sim a^2$ (or $ab$) (in case A of the purely transverse transport). It is possible that it is just the factor $R^{1.2} k^{0.5}$ that characterizes the quantitative ratio of the contributions of both types of transport to the energy balance for real tokamaks (the ‘dualism’ of energy transport).

Currently, the practical dualism of the local energy transport in tokamaks does not raise objections when it concerns either the SOL (Scrape-off layer) plasma, or the plasma with...
the destroyed magnetic surfaces in the regions of the tearing instability development (with characteristic plateaus at the $T_E(r)$ profiles) and in the regions of the developed sawtooth oscillations. That is, the dualism of the local energy transport is actually accepted. But the case is different, if we discuss the global energy transport. The ‘advocates’ of the transverse transport have a strong argument in favor of it: except for several ‘dead-end’ phenomena, such as the ‘isotopic effect’ or ‘anomalous pinch effect’, the observed ion transport is close in order of magnitude to that predicted by the neoclassical theory. For the majority of normal researchers, the global electron transport, which exceeds the neoclassical one by an order of magnitude or more, is the conventional neoclassical transport, but is enhanced due to the action of some instabilities, mostly of the gradient nature, which are many in the arsenal of theorists.

Experience teaches us that the transition from the conventional to the incomprehensible new follows a brutal algorithm, usually deeply painful. This experience can also be applied to formula (14). It was met at first as extremely heretical (although suitable for ‘interpolation estimates’). This dislike was extended to the ‘ITER scaling law’ (15), that follows from the fact that the ITER physicists gave him the condescending epithet ‘the engineering scaling law’ [27]. This label is a vivid illustration of the known maxim that ‘physicists, unlike engineers, begin to see a phenomenon only when they find an explanation for it’. As ‘an engineer’, Artsimovich ‘saw the phenomenon’. This indicated the anomalous nature of the electron losses in tokamaks, which first was announced by him in the aforementioned report in 1968 [6]. But the growth of the $\tau_E$ with plasma density (LOC and SOC modes, figure 5) put him as a ‘physicist’ at a dead end. I believe, he was the first to feel the dramatic nature of this challenge, but ideas of longitudinal electron transport [12, 13] were born only five years after his death.

I believe, in order not to conflict with his own and public principles, he reserved the right to recommend formula (14) ‘for the interpolation estimates for the steady state operating regime’ without dramatizing its global inconsistency with the fundamental principles of the magnetic confinement. Further, he accepted it as the ‘$\tau_E$ scaling law’ for the next step—the future T-10 tokamak. As it turned out later, his calculations were confirmed with high accuracy.

The drama of the $\tau_E$ behavior in relation to density was revealed later, when, against the background of the development of ideas for the practical use of the controlled fusion, it became clear that the thermal insulation of plasma ($\tau_E$) at all closed facilities (both at tokamaks and stellarators) steadily improves as its density increases. That is, it improves with an increase in the frequency of the Coulomb collisions. Besides, $\tau_E$ increases mostly in the major radius $R$, but not in the minor radius $a$. In all cases, the same anomalous electron thermal conduction becomes the key loss channel. For stellarators, this improvement could still be explained by the effect of the magnetic helical mirrors [28], but for tokamaks such an explanation is almost excluded. As is briefly outlined in the introduction, in an attempt to resolve this dramatic situation the concept of the longitudinal transport was developed [12, 13].

This concept seems to be quite suitable for the explanation of the following phenomena:

- the degradation of the energy lifetime $\tau_E$, which occurs with increasing electron heating power and becomes more pronounced, as the plasma is purified from impurities;
- a change in the sign of the electric potential of the core plasma from the negative to positive one;
- the ‘isotopic effect’;
- the ‘density pump-out’ effect;
- the ‘non-local transport’, which is characterized by the fast response (fractions of a millisecond) of the core plasma to the pulsed disturbances at its periphery;

and other less pronounced ‘inexplicable’ phenomena observed in tokamaks, like an SXR Pit. At the same time, the nature of the original mechanism of excitation of the magnetic flutter remains a mystery. Most often, opinions tend to be in the direction of ‘instability on trapped electrons’. On the other hand, the surprising similarity of energy scaling laws for closed systems—tokamaks and stellarators [28]—allows us to suspect that we may be talking about some unknown turbulences of a common element for them—Pfirsch—Schlutter currents. Moreover, similar magnetic disturbances are sometimes recorded on both devices by magnetic probes.

6. Conclusions

The SXR—Pit paradox is one of the phenomena that are not widely known. Within the framework of the concept of longitudinal transport, the above analysis made it possible to make the lower estimate of the length $L$ of the ‘assumed’ magnetic field line connecting the hot core of the plasma column with the ‘wall’ (the limiter) that is sufficient to ensure (within the framework of the classical model of the longitudinal electron heat conduction) the energy flows from plasma to the wall that are equal to the actually measured ones at the T-11M tokamak.

But under conditions of the limited length $L$ of the magnetic tube, the longitudinal transport should inherently ‘cut off’ the far tails of the Maxwellian electron energy distribution. This will occur as soon as the mean free path of the corresponding high-energy electrons in the surrounding ‘warm’ plasma with the characteristic time of collisions (13) [15] becomes equal to the length $L$ or exceeds it. It is just this ‘cutoff’ that manifested itself in the form of the SXR—Pit paradox, and was observed at the T-11M in the shots with low densities $N_e < 1.5 \times 10^{19} \text{ m}^{-3}$ (figures 2 and 3). Is it the ‘first barrier’ preventing the high-energy electrons of the Maxwellian distribution from transition to the runaway regime? This could become the subject of the subsequent studies of the physical nature of this phenomenon, which is critical for tokamaks.

Finally, the analysis presented made it possible to estimate the scale of the equivalent radial magnetic disturbance, $(b_r) \approx 0.1\% B_p$, which could cause the development of the longitudinal electron transport at the T-11M tokamak.

Taking into account that this magnetic disturbance is an order of magnitude lower than that characteristic of the tearing
instability, we can conclude that, at the T-11M tokamak, the energy store of the mechanism capable of transforming the transverse electron transport into the longitudinal one should be two orders of magnitude less than that maintaining the large-scale MHD activity in the plasma column. Thus, the search region for its physical nature is narrowed.

Summarizing the above, we conclude that, even without specification of its physical mechanism, the phenomenon of the SXR—Pit, obviously, associated with the ‘cutoff’ of the high-energy tails of the electron energy distribution, can actively motivate further studies of the physical nature of the anomalous electron transport in tokamaks. In this field of research, new discoveries can be expected by enthusiasts.

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Appendix. Lithium experiment at the T-11M tokamak at low and high plasma densities

The appearance of runaway electrons determines the lower limit of the electron density range for the tokamak operating regimes. The use of lithium coatings makes it possible to expand this density range due to the fact that the tokamaks with the lithium-coated walls can operate with the hydrogen recycling coefficient considerably lower than 1. The initial phase of the discharge in tokamaks characterized by the high loop voltage and low plasma density is the most dangerous by the generation of runaway electrons, since, in this phase, the initial group of runaway electrons can be generated, which then triggers the transition of all the electrons into the accelerating regime. Conventionally, the generation of runaway electrons (figure A1) is suppressed by increasing the working gas density in the initial phase of the discharge formation, which, at the hydrogen recycling coefficient close to 1, automatically sets the lower limit for the electron density \( n_e \) attainable during the entire subsequent discharge. However, if the recycling coefficient is lower than 1, it becomes possible to obtain plasma with the density \( n_e(t) \) decreasing with time to the specified value (figure A1 (B), curves 1, 2, and 3). This can be done by turning off the hydrogen inlet valve at some time of the discharge, and then the density \( n_e \) can be fixed at this level using the second specially selected valve.

In the lithium experiments at the T-11M with low densities, the initial peak of the electron density (\( n_e \)) averaged along the central observation chord was set at a level of approximately \( 2 \times 10^{19} \) m\(^{-3} \) (figure A1(B)) and, after turning off the valve, the density decreased to \( \sim 0.5 \times 10^{19} \) m\(^{-3} \) (figure A1(B), curve 1, 50 ms), or \( \sim 0.8 \times 10^{19} \) m\(^{-3} \) (figure A1(B), curve 2, 50 ms), or \( \sim 1.2 \times 10^{19} \) m\(^{-3} \) (figure A1(B), curve 3, 50 ms).

In the ‘high’—density experiments (figure A1(C)), the initial peak was set at a level of \( N_e \sim 4 \times 10^{19} \) m\(^{-3} \). As a result, it was possible to perform measurements in the density range from \( N_e < 0.5 \times 10^{15} \) m\(^{-3} \) to \( N_e > 4 \times 10^{19} \) m\(^{-3} \) without the occurrence of the runaway electron beams, as was confirmed by the signals of the control hard x-ray detector.

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