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On Measuring the Electron Heat Diffusion Coefficient in a Tokamak from Sawtooth Oscillation Observations

M. Soler
J. D. Callen

OAK RIDGE NATIONAL LABORATORY
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FUSION ENERGY DIVISION

ON MEASURING THE ELECTRON HEAT DIFFUSION COEFFICIENT
IN A TOKAMAK FROM SAWTOOTH OSCILLATION OBSERVATIONS*

M. Soler

Visitor from Junta de Energia Nuclear
Madrid, Spain

J. D. Callen

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ABSTRACT

A number of new methods are discussed for determining the electron heat conduction coefficient χ_e in a tokamak from the experimental observation of the space-time evolution of the temperature perturbations induced by internal disruptions. In the Oak Ridge Tokamak (ORMAK) the various average values of χ_e and the radial dependence of χ_e are found to be consistent with and more precise than the $\chi_e(r)$ determined by conventionally analyzing the electron power balance equation. The net result of these measurements is to conclusively prove that the dominant, radial electron heat transport mechanism in tokamaks is a microscopic, diffusive process.

1. INTRODUCTION

Since the dominant transport losses in a tokamak are apparently produced by radial electron heat transport, it is of fundamental importance to experimentally measure the radial electron heat conduction coefficient. The classical method of measuring a particular transport coefficient is to observe the system response to a small perturbation. Given the number of interrelated plasma parameters in a tokamak (density, electron and ion temperature, current density, impurity density, etc.), each governed by a separate transport equation, it would seem a nearly impossible task to measure a particular transport coefficient. Fortunately the naturally occurring internal disruptions in tokamaks [1, 2] cause small, cyclic temperature perturbations with little effect on other plasma parameters and thus provide an incisive probe for measuring the radial electron heat conduction coefficient χ_e in a tokamak [3].

In the first attempt [3] to measure χ_e by following the space-time evolution of the temperature perturbations induced by internal disruptions, the evolution was found to be diffusive but apparently governed by a coefficient about an order of magnitude larger than that governing the background plasma. This discrepancy was later resolved [2] through a better understanding of the internal disruption process which resulted in more correct initial conditions and "near-field" response considerations. It was concluded that the electron heat conduction coefficient is the same (within experimental error) as that governing the background plasma, and that this method represents a unique and very precise method of microscopically measuring the radial electron heat conduction coefficient.

In the process of a detailed reexamination of the heat pulse propagation problem [3] we have developed a wide variety of techniques for determining χ_e from the space-time evolution of the internal-disruption-induced temperature perturbations. In this work several methods are proposed for taking and analyzing the data to obtain both spatially averaged and local estimates of χ_e . The various methods are illustrated with ORMAK data of the same type and in many cases the same as that used in Refs [2] and [3]. With this set of albeit crude experimental data from ORMAK, determination of $\chi_e(r)$ from $r = 0$ to $r = 12$ cm is shown to be possible and entirely consistent with power balance inferences of $\chi_e(r)$ in ORMAK [4].

The paper is organized as follows. In the next section we discuss the way in which our understanding of the internal disruption process and experimental observations leads us to the initial conditions on the temperature perturbations. We solve the electron heat diffusion equation in Section 3 and show how the central (inside $q = 1$ surface) value of the heat conduction coefficient can be measured from the rate of the sawtooth rise. Next in Section 4 we discuss how the data can be analyzed to determine an average value of χ_e over the spatial range observed. Finally in Section 5 we discuss a method for providing a local measurement of χ_e , which requires more experimental information and accuracy. Some of the experimental techniques that should be used in taking the data are discussed in Section 6. We caution against sources of error in determining χ_e . Conclusions are summarized in Section 7.

2.. ELECTRON HEAT TRANSPORT PROCESSES IN THE COURSE OF INTERNAL DISRUPTIONS — GENERAL PHYSICAL DISCUSSION

The main features of electron heat transport taking place when internal disruptions appear are the following. At a certain instant of time during the discharge, the balance between sources of heat (ohmic heating, neutral injection, etc.) and sinks (radiation, internal conversion, etc.) which resulted in a certain steady temperature profile $T_c(r)$ for the plasma column, is abruptly flattened out to a certain radius, r_0 . Because of the formation of a magnetic island with a time scale much faster than the phenomena of normal heat diffusion, the electron heat diffusion coefficient $\chi_e(r)$ becomes practically infinite from $r = 0$ to $r = r_0$, where $r_0 \simeq \sqrt{2} r_s$, r_s being the singular surface $q = 1$. This results in a practically flat temperature profile throughout the mentioned range as shown in Fig. 1, so that the central temperature is lowered from $r = 0$ to $r = r_s$ and increased from $r = r_s$ to $r = r_0$ at a certain instant of time; the net or total heat content is unchanged (Fig. 1).

The resulting heat profile provides the initial condition for our problem with the points $r = r_s$ dividing regions of opposite phase in the heat perturbation. The phase starts downward inside $r = r_s$ and upward outside.

There is one particular point which we wish to discuss carefully because of its importance in setting up a mathematical model which is both simple and precise. After the disruption has taken place, thermal diffusion will tend to restore the system to a new equilibrium. The important point is whether this new equilibrium is reached and is the

same as the old one at the time the system engages in a new disruption. If not, the mathematical description becomes much more complex. We believe that with sufficient approximation, the old equilibrium is reached. To prove this point we observe that usually when the first disruption takes place during the discharge, the system is indeed in thermal equilibrium in the sense that no diffusive phenomena can be observed. After disruption, the soft x-ray pictures coming from the region $0 < r < r_s$ show a rise of the temperature stopping at practically the same previous level, with small differences changing randomly from pulse to pulse, before the next disruption takes place. After the disruptions stop, which happens several times in each shot in the ORMAK data we have analyzed, there is no sign that a significantly different equilibrium is reached. The electron temperature profile $T_e(r)$ should therefore be the same before, after, and at the end of each disruptive process.

These observations would seem contradicted by the somewhat puzzling fact that the temperature rise inside the $r = r_s$ surface looks linear when observed at different radii, and shows little or no sign of saturation. The explanation, we believe, lies mainly in the peculiar shape of these "internal sawteeth" which are shown as computationally calculated in Fig. 2. It can be seen that practically all the temperature rise comes from the "linear regime," with only 10-20% coming from the "saturation regime." The actual sawteeth observed may therefore show a rise of about 90% toward a new equilibrium temperature. This equilibrium would be about 10% higher than the former, and it is not surprising that possible changes in the local properties of the plasma, such as a momentarily lower impurity content, might account for this fact. In any

event, errors of about 10% in the description of this complex process seem acceptable. We will assume that ideally, in the absence of other complicated effects associated with changes in the nature of the plasma, the system returns to its former equilibrium state after each disruption.

3. MATHEMATICAL DESCRIPTION OF ELECTRON HEAT DIFFUSION AFTER INTERNAL DISRUPTIONS

3.1 BASIC EQUATION

It will now be shown that the simplification gained through the equilibrium assumption discussed above is considerable. The determination of the equilibrium state through a precise knowledge of the amount of radiation, current profile, etc., will no longer be our concern. Instead our subject will be the perturbation of this state with an equilibrium temperature profile $T_e(r)$ which can be experimentally observed.

Let us start with the general equation determining the electron temperature

$$\frac{3}{2} n_e \frac{\partial T_e(r,t)}{\partial t} = \frac{1}{r} \frac{\partial}{\partial r} n_e r \chi_e \frac{\partial T_e(r,t)}{\partial r} - \eta J^2 + \sum Q \quad (1)$$

where ηJ^2 represents ohmic heating, and $\sum Q$ represents other sources and sinks such as injection, radiation, recombination, etc. (We will suppress the subindex in T_e and n_e for simplicity.)

We will call $T_0(r)$ the equilibrium temperature profile just before disruption. Clearly it is determined by

$$\frac{1}{r} \frac{\partial}{\partial r} r n \chi \frac{\partial T_o}{\partial r} - n J^2 + \sum Q = 0 \quad (2)$$

Suppose now a disruption takes place at time $t = 0$, leaving the system in the state $T(r,0)$ (Fig. 1). The evolution from this state $T(r,0)$ at later times would be described by $T(r,t)$, which is governed by Eq. (1) subject to the initial condition $T(r,t) = T(r,0)$ at $t = 0$.

Rather than studying $T(r,t)$, we will study the evolution of the perturbation $\tilde{T}(r,t) = T(r,t) - T_o(r)$, which obviously satisfies the much simpler equation

$$\frac{3}{2} n \frac{\partial \tilde{T}(r,t)}{\partial t} = \frac{1}{r} \frac{\partial}{\partial r} r n \chi \frac{\partial \tilde{T}(r,t)}{\partial r} \quad (3)$$

and is subject to the initial condition

$$T(r,t) = \tilde{T}(r,0) = T(r,0) - T_o(r), \quad t = 0 \quad (4)$$

The shape of $\tilde{T}(r,0)$ is shown in Fig. 1.

Our task will be to extract as much information as possible about $\chi(r)$ from the experimental measurement of $\tilde{T}(r,t)$. Eventually of course this information should be fed into Eq. (1) so as to help determine the proper profiles of the sources and sinks.

3.2 EXPLICIT SOLUTION OF THE DIFFUSION EQUATION FOR THE PERTURBED TEMPERATURE

Before we try to determine $\chi(r)$, we should show that Eqs (3) and (4) indeed reproduce the experimentally observed diffusing temperature

perturbation profiles. In order to do this we will work out an approximate solution to Eqs (3) and (4), assuming that at least for radii not too large the density n and the coefficient χ can be taken as constant. Formally one has in this case

$$\tilde{T}(r, t) = \frac{a^2}{4\pi} \int_V \tilde{T}(r', 0) G_{t'=0} dV, \quad a^2 = \frac{3}{2\chi} \quad (5)$$

$G_{t'=0}$ represents the Green's function at time $t = 0$ in cylindrical coordinates, dV the volume element, and the integration is extended over the whole space, V . The general form of G is

$$G(\underline{r} - \underline{r}', t - t') = \frac{1}{t - t'} e^{-a^2 \underline{r} \underline{r}' \cos \theta / 4(t-t')} \theta(t - t'),$$

where $\theta(t - t')$ stands for the Heaviside step function, and $\underline{r}$, $\underline{r}'$ represent vectors. Since $(\underline{r} - \underline{r}')^2 = r^2 + r'^2 - 2 r r' \cos \theta$, the angular integration can be carried out in Eq. (5), according to

$$\int_0^{2\pi} d\theta e^{-a^2 r r' \cos \theta / 2t} = 2\pi I_0 \left(\frac{a^2 r r'}{2t} \right) \quad (6)$$

where I_0 is the modified Bessel function. So finally we have

$$\tilde{T}(r, t) = \frac{a^2}{2t} e^{-a^2 r^2 / 4t} \int_0^{r_0} r' dr' \tilde{T}(r', 0) e^{-a^2 r'^2 / 4t} I_0 \left(\frac{a^2 r r'}{2t} \right) \quad (7)$$

We have limited the integration at $r' = r_0$, since the perturbation $\tilde{T}(r, 0)$ as given by Eq. (4) vanishes outside this range. A reasonably good fit to the experimentally observed $T_0(r)$ at ORMAK is given by

$$T_o(r) = T_o \left[1 - \left(\frac{r}{r_e} \right)^2 \right]^2 \quad (8)$$

so

$$\tilde{T}(r,o) = T_o \left\{ \left[1 - \left(\frac{r}{r_e} \right)^2 \right]^2 - \left[1 - \left(\frac{r}{r_s} \right)^2 \right]^2 \right\} \quad (9)$$

Profiles for $T_o(r)$ different from Eq. (8) can be used, but the shapes of $\tilde{T}(r,t)$ are not sensitive to these changes. The shape given by Eq. (8) has the advantage of analytical simplicity.

The integrations involved in Eq. (7) cannot be carried out analytically. Numerical integration provides the result to any desired accuracy if we insert for $\tilde{T}(r',o)$ a series of δ functions in the form

$$\tilde{T}(r',o) \rightarrow \frac{r_o}{n} \sum_{i=1}^n \tilde{T}(r'_i,o) \delta(r' - r'_i) \quad (10)$$

where

$$r'_i = \frac{2i-1}{n} r_o$$

In practice we have seen that $n = 20$ is sufficient, and no appreciable changes appear for larger n . We usually "round off" $\tilde{T}(r',o)$ at $i = n$ making it two-thirds of its calculated value and add the remaining one-third as a new $(n+1)$ term to the series in Eq. (10). An insignificant error is thus made, and this provides a simple way to imitate rounding off the physical $\tilde{T}(r,o)$. Different curves corresponding to different radii r are shown in Fig. 3. Several features of these curves are of

practical importance since they lend themselves to physical observation. First, their peak tends to be delayed as we go to larger r . Second, they become negative at a certain time t_z , which also increases for larger r . Third, the negative tail also goes through a minimum. Fourth, the ratio between $\tilde{T}$ at the maximum and at the minimum decreases with r (in other words, the minimum becomes relatively deeper with increasing radius). All of these features are clearly observed in our set of data, so we are confident that even with the approximations taken, our model is correct.

3.3 DETERMINATION OF $\chi(o)$

Equation (6) clearly simplifies when $r = 0$. We will therefore try to evaluate χ at this point. Since

$$I_o \left(\frac{a^2 r r'}{2t} \right) = e^{-a^2 r^2 / 4t} = 1, \quad (r = 0)$$

we have

$$\tilde{T}(o, t) = \frac{a^2}{2t} \int_0^{r_o} r' dr' T(r', o) e^{-a^2 r'^2 / 4t} \quad (11)$$

For $\tilde{T}(r', o)$ given by Eq. (8), Eq. (9) can be easily integrated. If we write $\tilde{T}(r', o) = A + Bx + Cx^2$ with $x = r'^2$ and call $a^2/4t = \alpha$, we have for small t

$$\frac{\partial \tilde{T}(o, t)}{\partial t} \simeq \frac{4}{a^2} \left[B - (B + 2r_o^2 C) e^{-\alpha r_o^2} \right] \quad (12)$$

Since for Eq. (8), $B = 2/r_\ell^2 T_o$, $C = -T_o/r_\ell^4$,

$$\frac{1}{T_o} \frac{\partial \tilde{T}(o,t)}{\partial t} = \frac{16\chi}{3r_\ell^2} \quad (13)$$

For profiles different from Eq. (8), the term with the exponential in Eq. (12) in general will not vanish, but will be negligible, so Eq. (13) is quite exact.

Experimentally, for our set of data (ORMAK discharges 16193-7) $T(o,t)$ rises approximately linearly (as was explained in Section 2) to its former level in about 2 msec. Since $r_\ell \approx 20$ cm and $\Delta T/T_o \approx 20\%$, we find

$$\frac{1}{5.2 \times 10^3} = \frac{16\chi(o)}{1200}, \quad \text{or} \quad \chi(o) \approx 0.8 \times 10^4 \text{ cm}^2/\text{sec}$$

We have written $\chi(o)$ and not just χ because the only contribution to the integral in Eq. (11) comes from the origin, where the assumptions of constant χ and n amount to no restriction. In other words, the diffusive process at $r = 0$ is almost a strictly local phenomenon and should not be affected by variations in n which are only significant at much larger radii. The same result can be obtained for $t = 0$ from the simple use of Eq. (8) in Eq. (3). Equation (14) probably provides the simplest measure of $\chi(o)$ one can get in a tokamak plasma.

At this point it may be useful to explain that our assumptions at $r = 0$ are equivalent to those made in Ref. 2. In order to see this clearly we may write Eq. (1) in the form

$$\frac{\partial T}{\partial t} = \frac{\partial \tilde{T}}{\partial t} = \frac{2}{3} \frac{1}{n} \left(\frac{1}{r} \frac{\partial}{\partial r} r \chi_e T_o + \frac{1}{r} \frac{\partial}{\partial r} \chi_e \tilde{T} + \underline{\eta J_z^2 - \Sigma Q} \right)$$

where we have underlined differently those terms which may by themselves constitute two separate equations. Since the shapes of T_o and $\tilde{T}$ are the same in the origin, with only opposite signs as shown in Fig. 1, clearly one can write for $r = 0$

$$\frac{\partial T}{\partial t} = \frac{\partial \tilde{T}}{\partial t} = \frac{2}{3} \frac{1}{n} (\eta J_z^2 - \Sigma Q)$$

This is Eq. (3) in Ref. 2. It was also argued in that reference that the term ηJ_z^2 would dominate over $-\Sigma Q$, thus explaining the linear rise of T . We have discussed in Section 2 the limitations of this argument.

4. DETERMINATION OF AVERAGE VALUES OF χ UP TO A CERTAIN RADIUS AND RESOLUTION OF THE ORIGINAL HEAT PULSE PROPAGATION PARADOX

4.1 COMPARISON OF HEAT PULSES WITH COMPUTED $\tilde{T}(r,t)$

We will now obtain values of χ away from the origin but for small enough r to justify our assumptions of constant n and χ . The simplest procedure is to fit our theoretical curves $\tilde{T}(r,t)$ at a certain r and adjust the value of χ until the experimental curves are matched. The smaller the r , the more reliable the fit will be. In practice this means taking the r corresponding to the first experimental chord outside r_o . The fit is simple to carry out since one needs only take χ such that t_z [the distance from the origin to the zero of $\tilde{T}(r,t)$] equals the experimental value. Within the experimental uncertainties, the theoretical $\tilde{T}(r,t)$ fits well the observed heat pulse for $\chi = 1.8 \times 10^4$ cm²/sec. This should represent the average value $\bar{\chi}$ from $0 \leq r < 8$ cm.

4.2 DESCRIPTION OF A SIMPLE DIPOLE MODEL

Although the general solution $\tilde{T}(r,t)$ cannot be given in a closed analytical form, a rough but very useful analytic approximation is obtained by making $n = 2$ in Eq. (10). This amounts to calculating the heat diffusion due to two radially localized sources of opposite sign and zero total heat content. We will call this initial condition a simple dipole source. A large n in Eq. (10) implies taking a superposition of simple dipoles located at different points, having the effect of an extended dipole source with similar properties, particularly at large distances.

There is a great deal of freedom in constructing a simple dipole. One can easily realize that Eq. (10) cannot be blindly applied for small values of n . For $n = 2$ the proper procedure is to locate the two δ -like contributions, respectively, at some point in the intervals $0 < r_1 < r_s$ and $r_s < r_2 < r_o$, so that the total heat content is zero, that is,

$$\int_0^{r_o} r \tilde{T}(r,0) dr = r_1 r_s T_- + (r_o - r_s) r_2 T_+ \quad (14)$$

Equation (14) still leaves some arbitrariness which we will later adjust.

We can now write Eq. (6) in the form

$$\begin{aligned} \tilde{T}(r,t) \approx & \frac{a^2}{2t} e^{-a^2 r^2/4t} \left[r_1 r_s T_- e^{-a^2 r_1^2/4t} I_0\left(\frac{a^2 r r_1}{2t}\right) \right. \\ & \left. + r_2 (r_o - r_s) T_+ e^{-a^2 r_2^2/4t} I_0\left(\frac{a^2 r r_2}{2t}\right) \right] \end{aligned} \quad (15)$$

or using Eq. (11)

$$\tilde{T}(r,t) \approx \frac{a^2 A}{2t} e^{-a^2 r^2/4t} \left[e^{-a^2 r_1^2/4t} I_0\left(\frac{a^2 r r_1}{2t}\right) - e^{-a^2 r_2^2/4t} I_0\left(\frac{a^2 r r_2}{2t}\right) \right], \quad A = r_1 r_s \dot{T}_- \quad (16)$$

There are two well-known expansions of $I_0(z)$ which help in discussing Eq. (16). For large and small arguments one has, respectively

$$I_0(z) \approx \frac{e^z}{\sqrt{2\pi z}} \left(1 + \frac{1}{8z} + \dots \right) \quad (17)$$

$$I_0(z) \approx 1 + \frac{\left(\frac{1}{4} z^2\right)}{(1!)^2} + \frac{\left(\frac{1}{4} z^2\right)^2}{(2!)^2} + \dots$$

Equation (16) with the help of Eq. (17) allows one to deduce with very little effort the shapes of the heat pulses $\tilde{T}(r,t)$ at any r and t . They differ little from the numerically calculated ones (e.g., $n = 20$) and fit well the experimental data.

4.3 A SIMPLE FORMULA FOR $\bar{\chi}$ UP TO A CERTAIN RADIUS

From the experimental point of view one of the simplest observables that can be measured in the external (or $r > r_0$) heat pulses is the width of the pulse from the origin to the point where it becomes negative. This is given by Eq. (16) as the solution of

$$e^{-a^2 r_1^2/4t} I_0 \left(\frac{a^2 r r_1}{2t} \right) - e^{-a^2 r_2^2/4t} \left(I_0 \frac{a^2 r r_2}{2t} \right) = 0 \quad (18)$$

Using the expansion of I_0 for the case of large parameter we have

$$e^{-a^2 r_1^2/4t} e^{-a^2 r r_1/2t} = \sqrt{\frac{r_1}{r_2}} e^{-a^2 r_2^2/4t} e^{-a^2 r r_2/2t}$$

or

$$\frac{2a^2 r(r_1 - r_2) - a^2(r_1^2 - r_2^2)}{4t_z} = \ln \sqrt{\frac{r_1}{r_2}}$$

For simplicity we may use the reasonably good approximation $r_1 + r_2 = 2r_s$ for the location of the dipole sources, so we finally obtain

$$\chi \approx \frac{3(r_1 - r_2)(r - r_s)}{2t_z \ln(r_1/r_2)} \quad (19)$$

We must now use our freedom in Eq. (14) and locate properly the dipole sources to best represent the extended dipole. In fact one finds that for plausible choices of r_1 and r_2 the numerical factor

$$r_1 - r_2 / \ln(r_1/r_2)$$

is rather constant. A reasonable choice is $r_2 = 1.3r_s$ and $r_1 = 0.6r_s$, which gives a numerical factor very close to unity. For $r_s = 5.5$ and $r = r_0 = \sqrt{2} r_s$ (i.e., assuming r is located at the first observable chord outside r_s) we find

$$\bar{\chi} = \frac{1.8 \times 10^4}{t_z} \text{ cm}^2/\text{msec}$$

For $t_z = 0.8 \text{ ms}$, $\bar{\chi} \approx 2.2 \times 10^4 \text{ cm}^2/\text{sec}$, the agreement with the numerical fit is very good.

Equation (19) is quite interesting, not only because it makes almost unnecessary the recourse to numerical calculation, but because it relates in a very simple way $\bar{\chi}$ with a quantity, t_z , that can be observed throughout a discharge. As we will see, these average values of χ are subject to rather large errors because of the assumption of constant n , so one should readjust the numerical factor $(r_1 - r_2)(r - r_g)/l(r_1/r_2)$ corresponding to a certain r . Once this is done, if we keep at the same r and observe a sequence of sawteeth throughout the discharge, we obtain directly the time variation of $\bar{\chi}$ by just measuring t_z .

Experimentally it is known that the energy containment time increases during the discharge. So does t_z , by about 50%, so we can state that χ decreases by this amount from 15 to 60 msec in UKMAK.

4.4 OTHER DETERMINATIONS OF $\bar{\chi}$ OVER EXTENDED REGIONS AND RESOLUTION OF A PARADOX

A brief description will now be given of a method for measuring $\bar{\chi}$ based on observation of the time of arrival t_p of the peak of the heat pulses as they are observed from different radii. This method has historical interest since it was the first to be used for determining χ from heat pulse propagation [2, 3]. For reasons that will become clear, the χ value obtained appeared much higher than expected from rough energy containment time calculations.

If we look at the behavior of $\tilde{T}(r,t)$ in our simple dipole approximation Eq. (16) at large t , we may approximate the Bessel function by the small parameter expansion. This gives

$$\begin{aligned}\tilde{T}(r,t) &\sim \frac{a^2 A}{2t} e^{-a^2 r^2/4t} \left(1 + \frac{1}{4} \frac{a^4 r^2 r_1^2}{4t^2} - 1 - \frac{1}{4} \frac{a^4 r^2 r_2^2}{4t^2} + \dots \right) \\ &= \frac{a^6 A r^2}{32} \frac{1}{t^3} e^{-a^2 r^2/4t}\end{aligned}\quad (20)$$

The maximum of $\tilde{T}(r,t)$ lies at

$$\frac{\partial \tilde{T}}{\partial t} \sim \left(-\frac{3}{t^4} + \frac{a^2 r^2}{4t^5} \right) e^{-a^2 r^2/4t} = 0$$

so

$$t_p = \frac{a^2 r^2}{12} = \frac{r^2}{8\chi} \quad \text{or} \quad \chi = \frac{r^2}{8t_p} \quad (21)$$

If we plot r^2 vs. t_p , the observed slope should be a straight line giving directly χ for large values of r^2 . In the first attempt at measuring χ by this method, the initial condition or heat pulse disturbance at $t = 0$ was approximated by just one δ function. Instead of Eq. (17) one had

$$\tilde{T}(r,t) \sim \frac{e^{-a^2 r^2/4t}}{t},$$

which peaks at

$$\frac{\partial \tilde{T}}{\partial t} \sim \left(-\frac{1}{t^2} + \frac{a^2 r^2}{4t^3} \right) e^{-a^2 r^2/4t} = 0 \quad .$$

so

$$t_p = \frac{a^2 r^2}{4} \quad , \quad \chi = \frac{3r^2}{8t_p}$$

a factor of three higher than Eq. (21). These asymptotic arguments are not completely precise, although they give a clear idea of the origin of the paradox. In fact, the exponentials in Eq. (16) should also be expanded, which gives instead of Eq. (18) a value about 10% lower in the first approximation. But also the expansion of the Bessel function should be changed to that for large arguments at small r (r close to r_0). A set of curves $t_p = t_p(\chi, r_p^2)$ is thus obtained which must be compared with the experimental data (Fig. 4). The results obtained for a certain set of data from ORMAK are shown in Fig. 5. The great discrepancy found in Ref. [3] was then reduced to an acceptable level given the approximations made and the experimental uncertainties involved. The value obtained, $\bar{\chi} = 3.5 \times 10^4 \text{ cm}^2/\text{sec}$, is an average over the region, $0 \leq r < 14 \text{ cm}$, where the errors coming from the neglect of $\partial n/\partial r$ are probably very significant. Other important sources of error are related to the difficulty of defining the peak of the heat pulse. The time delays are extremely small and the relative errors consequently can be very large. This is complicated more by the existence of a great deal of $m = 1$ magnetic activity at the moment of disruption [5], which coincides with the origin of the heat pulse and frequently hides it completely. Another drawback of the method is its essentially nonlocal

or average nature. In order to fit a curve, $t_p = t_p(r^2)$, we need a number of points at different r extending to a region of about two-thirds of the plasma column where the approximation of constant n is quite incorrect. It would seem therefore that the other methods discussed in this paper should be used in preference to this method.

5. EXACT CALCULATION OF LOCAL $\chi(r)$ FROM HEAT PULSE PROPAGATION

The value $\chi \approx 0.8 \times 10^4$ obtained for $r = 0$, compared with the average values $\bar{\chi} \approx 1.8 \times 10^4$ for $0 \leq r < 8$ and $\bar{\chi} \approx 3.5 \times 10^4$ for $0 < r < 14$, seems to indicate that the effect of our main approximation $n \approx \text{constant}$ for obtaining the averages is to considerably increase $\bar{\chi}$ for large radii. In fact this is not surprising. If we allow an r dependence in n , a new term should be added in Eq. (3)

$$\frac{3}{2} \frac{\partial \tilde{T}}{\partial t} = \frac{1}{n} \frac{\partial n}{\partial r} \frac{\partial \tilde{T}}{\partial r} \chi + \frac{1}{r} \frac{\partial}{\partial r} r \chi \frac{\partial \tilde{T}}{\partial r} \quad (22)$$

The derivatives $\partial \tilde{T} / \partial r$ are negative for most of the values of r and so is $\partial n / \partial r$. The average values of $\bar{\chi}$ obtained by neglecting the term with $\partial n / \partial r$ in Eq. (19) should therefore be higher than in reality.

This observation leads to the attractive possibility of measuring χ precisely at those points where $\partial \tilde{T} / \partial r = 0$ and where the assumptions of constant n and χ would not be necessary. We thus obtain local values of χ exact even for variable $n(r)$

$$\chi(r) = \frac{3}{2} \frac{\partial \tilde{T} / \partial t}{(\partial^2 \tilde{T} / \partial r^2)}, \quad \text{if} \quad \frac{\partial \tilde{T}}{\partial r} = 0 \quad (23)$$

This formula is only subject to the very general assumptions involved in arriving at Eq. (3).

In order to calculate $\chi(r)$ accurately from Eq. (23), it would be convenient to have a sufficient high number i of chords measuring $\tilde{T}(r_i, t)$ so that the whole surface $\tilde{T}(r, t)$ might be accurately reconstructed by interpolation of curves of constant t . We would then look for those points where $\partial\tilde{T}/\partial r = 0$ and measure at those points the slope $\partial\tilde{T}/\partial t$ and the curvature $\partial^2\tilde{T}/\partial r^2$. In general, there is no reason why these quantities should be small or hard to measure, so the resulting $\chi(r)$ should also be accurate.

The accuracy to be expected from this method can easily be checked in the case of constant n and χ . We can calculate the theoretical curves, $\tilde{T}(r_i, t)$, at different r_i and interpolate the surface for each constant t . Assuming, as is the case with ORMAK, that there are four detectors falling outside $r = r_0$ (at $r_{1,2,3,4} = 8, 10, 12, 14$ cm), we have run a program for $r_s = 5.5$ and $\chi = 1.5 \times 10^4$ cm²/sec to obtain $T(r_i, t)$ and then determined by Lagrange interpolation the whole grid $T(r, t)$. We have then recalculated $\chi(r)$ using Eq. (23). The results are given in Table I. Except at the end points, where the errors of interpolation are large for the second derivative, the results are quite good. There is no reason why the errors in the variable n case should exceed those of Table I. By increasing n to 6 or 8, the accuracy in calculating $\chi(r)$ would be as exacting as can be desired.

As a crude example of the method, we show (in Fig. 6) the average shapes of the heat pulses at four radii in ORMAK coming from five consecutive discharges (No. 16193-7) which had very similar features. The

TABLE I. VALUES OF $\chi_e(r)$ AS RECALCULATED FOR THE CASE OF CONSTANT χ_e

Time (msec)	Radius (cm)	χ_e ($\times 10^4$ cm ² /sec)
0.46	8.50	2.059
0.54	9.00	1.733
0.66	9.50	1.564
0.80	10.00	1.506
0.94	10.50	1.502
1.08	11.00	1.511
1.22	11.50	1.514
1.36	12.00	1.496
1.68	13.00	1.355
1.86	13.50	1.237

internal disruption selected occurred at practically the same time, about 35 msec. The surface $\tilde{T}(r,t)$ is traced by interpolation in perspective, and the proper shapes of different curves at constant times are drawn in Fig. 7. Table II shows the measurements taken directly from the curves.

In spite of the crudeness of our calculations (Fig. 8) there is a very striking agreement between $\chi(r)$ obtained in Table II, $\chi(0)$ obtained in Section 3, the average results $\bar{\chi}$ obtained assuming constant n , which comes out too high as we had predicted, and the $\chi_e(r)$ inferred from unfolding the radial electron heat balance [4].

6. DISCUSSION OF THE EXPERIMENTAL MEASUREMENTS

One clear advantage of using Eq. (3) is that relative increments $\tilde{T}$ are much easier to measure accurately than absolute temperatures T by soft x-ray techniques. However, there are some problems such as the effect of filters which must be dealt with properly in order to avoid errors in the method advocated in Section 5. It should be realized that the relative amplitudes of $\tilde{T}(r_i, t)$ at different r_i must be known in order to calculate $\chi(r)$ as in Eq. (24). This would not be necessary for calculating $\bar{\chi}$ as explained in Section 4. The possible effect of impurities at different r_i should perhaps also be estimated. Clearly, the method of electron cyclotron resonance [6] would be best suited for calculating $\chi(r)$ if profiles coming from enough radii were available.

It should be stressed that direct x-ray pictures or information, as opposed to information coming from filtered signals, is essential. In ORMAK the detailed information coming through a high-pass filter proved

TABLE II. ESTIMATES OF χ USING MEASUREMENTS FROM FIGS. 6 AND 7

Points	Δ (cm)	$\tilde{T}_i$	$\tilde{T}_2$	$\tilde{T}_o$	$\frac{\partial^2 \tilde{T}}{\partial r^2}$	$\frac{\partial \tilde{T}}{\partial t}$	χ
1	1	5.6	5.6	6.7	2.2	16	1.1×10^4 cm ² /sec
2		3.0	3.0	3.5	1.2	10	1.2×10^4 cm ² /sec
3		1.4	1.4	1.8	0.8	5.8	1.1×10^4 cm ² /sec
4		0.5	0.5	0.65	0.3	3.6	1.8×10^4 cm ² /sec
5		0.1	0.1	0.15	1.0	0.1	1.5×10^4 cm ² /sec

useless. This is because the slopes $\partial\tilde{T}/\partial t$ are drastically changed except at the peak, and in trying to reconstruct the original pulses, the small differences between filters corresponding to each channel introduced enough distortion in $\tilde{T}(r,t)$ to make the calculation in Eq. (23) irrelevant.

7. SUMMARY

A general analysis has been made of the properties of the surface $\tilde{T}(r,t)$ describing the electron heat temperature perturbations produced by internal disruptions in ORMAK. Several observable quantities have been shown to provide measurements of χ_e , the electron heat diffusion coefficient. In particular, χ_e can be calculated exactly at $r = 0$ [cf., Eq. (13)] and, from different observables, at those points in which $\partial\tilde{T}/\partial r = 0$ [cf., Eq. (23)]. Approximate average estimates $\bar{\chi}_e(r_i)$ from the origin to some radius r_i can also be made on the assumption of constant density [cf., Eqs (19), (21)]. However, since the plasma density decreases with radius, these should give a higher value of χ_e and should be less reliable as r_i becomes larger. Some of these average measurements are extremely simple to make and have the advantage of providing accurate information on the evolution of χ_e as the discharge advances in time.

With the set of data from ORMAK which we have analyzed, a consistent picture emerges. It is compatible with other volume-average calculations made to determine χ_e [4]. Although our measurements are taken over an unexpanded x-ray oscilloscope picture having poor resolution, our estimated error bars give an idea of the accuracy of the method.

Careful experimental measurements extending all the way out to the plasma edge should be undertaken in order to determine $\chi(r)$ by this method with the desired accuracy. In ORMAK all the discharges which exhibit internal disruptions seem to have roughly the same energy confinement time and hence χ_e . Thus, we have not been able to perform any scaling checks of how χ_e determined from the space-time evolution of the temperature perturbations compares with that determined from the gross energy confinement time. Hopefully this will be done in future experiments.

Finally, we would like to reiterate that the most fundamental conclusion of this work is that between disruptions the temperature perturbations produced by internal disruptions evolve by a diffusive process, at least on length scales longer than about 1.7 cm in ORMAK — the scale of spatial resolution in the soft x-ray system. In particular, the transport does not appear to be wavelike or ballistic in nature as might occur with macroscopic flows in a plasma. In this paper we have demonstrated through a variety of analytic methods of the experimental data that the electron heat conduction coefficient governing the temperature perturbations is the same, to within experimental error, as that governing the background plasma. In this regard it is important to note [2] that internal disruptions can occur only in plasmas for which the impurity radiation from the center of the plasma is less than local ohmic heating power input. Thus, the data discussed herein are undoubtedly characteristic of ORMAK discharges where the dominant heat loss mechanism is by electron heat transport and not by impurity radiation. The net result of these observations is to prove that when

electron heat transport is dominant, it is a microscopic diffusive loss process. Since the favorable scaling of confinement time with plasma size ($\tau_E \sim a^2/4\chi$) is based on the premise of a diffusive loss process, these observations bode well for the future of tokamaks as long as the diffusion coefficient can be kept small.

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FIGURE CAPTIONS

FIG. 1. Internal disruptions perturb the temperature equilibrium profile T_e . The resulting temperature profiles are flat from the center to a certain radius r_0 . We call r_s the radius of disruption, where the temperature should show no change. For the particular profile shown, which imitates reasonably well the profiles found at ORMAK, $r_0 \approx \sqrt{2}r_s$. Typically r_s is contained within the limits shown in the figure. The hatched profiles $\tilde{T}_1(r,0)$, $\tilde{T}_2(r,0)$ represent the corresponding initial conditions for the perturbation Eq. (3).

FIG. 2. Typical shape of internal sawteeth, as calculated from a transport code. The saturation part of the curve is seldom seen in experiments in ORMAK, but the temperature rise in most cases approximately equals the temperature before disruption, randomly lying slightly above or below it.

FIG. 3. Different shapes of $\tilde{T}(r,t)$ as theoretically calculated for constant n and χ . For generality, dimensionless units are used: $R_i = r/r_s$, $\tau = a^2 t$ with $a^2 = 3/2$.

FIG. 4. Using again dimensionless units, general curves $\tau = \tau(r^2)$ are plotted for different initial conditions, corresponding to the extreme values of r_s in Fig. 1. These curves must be used to fit the experimental time delay of arrival of the heat pulse peak at different radii.

FIG. 5. Distribution of experimental results for χ from heat pulse propagation (as measured from peak time of arrival) χ_{ep} , vs the gross energy containment time estimate of χ_{ec} for the same discharge.

FIG. 6. Interpolation of the surface $\tilde{T}(r,t)$ by means of different lines of constant time t . The four curves $\tilde{T}(r_i,t)$, $r_i = 8, 10, 12$, and 14 are drawn by measuring in photographic ORMAK data the absolute maximum, the first zero t_z , the location of the minimum, and following the general shape of the curves as calculated in the case of constant n and χ .

FIG. 7. Actual shape of the curves of constant t , in $\tilde{T}(r,t)$. The second derivative $\tilde{T}'' = \partial^2 \tilde{T} / \partial r^2$ is calculated at the points of $\partial \tilde{T} / \partial r = 0$ using the parabolic formula $T'' = (T_1 - 2T_0 + T_2) / \Delta^2$ directly in the drawing with $\Delta r = 1$ cm; the results are shown in Table II.

FIG. 8. Determination of $\chi_e(r)$ in ORMAK (discharges 16193-7) using Eq. (13) and results from Table II. We have included an approximate estimate of the errors for the three measurements. The dotted lines represent the two upper bounds corresponding to an average radii up to $r = 8$ and $r = 14$. We also include the results of Ref. [4], Fig. 11 corresponding to inclusion (lower continuous line) and exclusion (upper continuous line) of radiation.

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Fig. 1.

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$$\frac{\Delta \tilde{T}_{\text{saturation regime}}}{\Delta \tilde{T}_{\text{linear regime}}} = 11\%$$

Fig. 2.

Fig. 3.

Fig. 4.

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Fig. 5.

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Fig. 6.

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Fig. 7.

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Fig. 8.

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