HEATING A TOKAMAK PLASMA

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Heating a Tokamak Plasma*

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ABSTRACT

An attempt is made to assess the role of heating in the context of the entire problem of the attainment of controlled fusion. Attention is directed to toroidal configurations — primarily tokamaks — and we examine in turn the dominant processes for the transport of heat from the plasma across the magnetic field to the wall: thermal conduction, bremsstrahlung, and synchrotron radiation. Balancing this heat loss is the input from Ohmic heating, alpha particles, and from supplementary sources. A zero-dimensional neoclassical model is used to find asymptotic conditions. This simple model enables us to make a quick survey over a wide range of device parameters and to evaluate the need for and the effectiveness of given amounts of supplementary heating power. It is found that really modest amounts of added heat can introduce impressive changes in laboratory plasmas, and can greatly simplify the ignition of a reactor plasma.

INTRODUCTION

As an introduction to this Symposium on Heating and Injection, it seemed to me appropriate to examine the role of heating in the overall framework of our attempt to achieve controlled fusion. I will direct my remarks almost exclusively to toroidal devices—particularly tokamaks—and the two major questions I will try to answer are, first, "Is Ohmic heating by itself sufficient to achieve ignition in a fusion reactor?" and second, "What are the effects on the plasma due to the addition of a given amount of supplementary power?" I will not try to differentiate between the different means for supplying the additional heating power, such as rf heating, injection, compression, shocks, turbulence, lasers, etc. We will hear specific details about these various methods over the coming days and the qualitative differences between them should become quite clear. On the other hand, we may, in simplest terms, regard each of these heating schemes as a means to deposit so many watts or joules in the plasma. Without worrying about whether the heating power goes initially into the ions or the electrons, into the edge of the plasma or into its center, and without inquiring in more than a cursory manner about instabilities, convection, and enhanced loss which might accompany the heating process, we will try to look at the major classical heat transport processes which will compete against any heating system. These are three in number: thermal conductivity, bremsstrahlung, and synchrotron radiation. On the other side of the power balance equation, in addition to
supplementary heating, we will compute the amount of plasma heating due to neoclassical Ohmic heating and also due to alpha particles from the D-T nuclear reaction.

**Thermal Conductivity.** For our present experimental plasmas, the dominant mechanism for heat loss is thermal conductivity. Nevertheless, perhaps the first question one might raise is whether one could design a reactor that had only a negligible amount of thermal conduction loss from the plasma. For example, could one maintain constant temperature all the way out to the edge of the plasma? The answer is not necessarily a categorical no. For example, in the conceptual design of mirror-geometry fusion reactors, it is generally assumed that particles leaving the plasma are deposited in a perfect sink, and that there is no recycling. If one could produce a divertor for a tokamak which had these properties, then heat loss would be much reduced, the stability of the plasma would be improved, and the problem which faces us at this Symposium — that of heating the plasma — would be considerably simpler. On the other hand, while mirror engineers are willing to devote much more space to their divertor than to the mirror-machine itself, nevertheless the toroidal field design forces tokamak engineers to be very stingy with the allotment of space to a divertor. Figure 1 shows a tokamak reactor design which is currently under study at the Plasma Physics Laboratory, in Princeton. It has a strip divertor of considerable area which goes all the way around the torus. Nevertheless, the density of ions flowing into the divertor would be of the order of $10^{12} \text{cm}^{-3}$, and with a plasma of this density in the divertor
channel, there would seem to be no way to prevent cold electrons released at the target wall from streaming back to the edge of the plasma.

Based on our present ideas, therefore, we must assume that the edge of the toroidal plasma will be quite cold, and we therefore turn now to look at the mechanism of thermal conduction across a magnetic field.

With classical processes — that is, on the basis of two-body Coulomb collisions — one expects that it will be the ions that transport most of the heat across the magnetic field. That is simply because the step length, which is the Larmor radius, is much larger for ions than for electrons. It is for a different reason that ion heat transport across the magnetic field goes at a faster rate than ion particle transport. When two particles of like charge and mass collide, their individual energies may be redistributed, but their center of gravity does not move. Moreover, transport of ions and electrons across the magnetic field must proceed at the same rate, and it turns out that this rate is down by the square root of the mass ratio from the rate of heat transport. Therefore, if we direct our attention only to classical heat transport processes, we can neglect the heat which is carried out bodily with each particle.

One can make a qualitative estimate of the magnitude of the coefficient for thermal conductivity. We look for a heat flow equation of the type

\[ \dot{Q}_\perp = -\kappa_\perp \nabla k T \]
\[ k_\perp = n_\perp D_\perp, \]
\[ D_\perp \sim \frac{\lambda^2}{\tau} + \frac{\rho_{Li}^2}{\tau_{90^\circ}}. \]

Here \( D_\perp \) represents a diffusion coefficient for transport across the magnetic field, \( \rho_{Li} \) is the ion Larmor radius in the toroidal field, and \( \tau_{90^\circ} \) is the 90° deflection time for ion-ion collisions.

Now, in the long mean-free-path regime, there is a step length longer than \( \rho_{Li} \) which must be used. That is the thickness of the ion banana orbits. Figure 2 shows a sketch of the orbits in a tokamak. They may be divided into two categories — passing particles and trapped particles. The trapped particles have velocity vectors close to 90° from the magnetic field direction, and are reflected by the somewhat stronger magnetic field on the inside of the torus. The total thickness of the banana orbit at its fattest point is just \( 2^{3/2} \) times the Larmor radius in the local \textit{poloidal} field, \( B_{pol}(r) \), divided by the square root of the local aspect ratio \( (R/r) \). We can now make a new estimate of the heat conductivity in what is called the neoclassical or banana regime. To do this, however, we need first to look briefly at physics of magnetic mirror reflection. Consider a particle with magnetic moment \( \mu = mv_\perp^2/2B \) situated between two magnetic mirrors. At the point of reflection (designate by \( B = B_{refl} \)), the parallel velocity will go to zero, and \( v_\perp^2 \) at this point will be equal to \( v^2 \), where \( mv^2/2 \) is the total particle kinetic energy. Now going back to the original particle location in between the mirrors,
we can write down the condition for reflection, or for trapping,

\[
\frac{mv^2}{2B} = \frac{mv^2}{2B_{\text{refl}}}
\]

\[
\sin^2 \theta \frac{v^2}{v^2} = \frac{B}{B_{\text{refl}}} \frac{R-r}{R} = 1 - \frac{r}{R}
\]

\[
\cos \theta = \sin \left( \frac{\pi}{2} - \theta \right) \approx \sqrt{\frac{R}{r}}
\]

here \( R \) is the major radius of the torus, and \( r \) is the minor radius coordinate for the particle.

Now we have all the information necessary to estimate the neoclassical heat conductivity. The heat is carried mainly by the trapped ions — therefore we reduce \( n_i \) by the ratio of the trapped-particle solid angle to \( 4\pi \), i.e., by \( \int_0^{\pi/2} d(\cos \theta) = \cos \theta \approx (r/R)^{1/2} \). Also we reduce the angle of deflection from 90° to \((\pi/2 - \theta)\), which is the critical angle for scattering into the loss cone. This effect reduces \( \tau_{90^\circ} \) to \( r/R \) times its earlier value. Putting it together, we have

\[
\bar{Q}_i \approx - n_i \frac{(r/R)^{1/2}}{B_{\text{pol}}} \left( \frac{\rho_{\text{Li}} B}{B_{\text{pol}}} \frac{(r/R)^{1/2}}{2} \right)^2 \tau_{90^\circ}(r/R) \kappa T
\]

The actual coefficient, determined by a sophisticated mathematical analysis,\(^1\) leads to the form for neoclassical heat conduction which we will use,
\[ P_Q = 2.17 \times 10^{-22} \frac{Z^2 \text{eff} n_i^2 A^{1/2}}{B^2 R^2} q^2 T_i^{1/2} \ln \Lambda \frac{B}{R} \frac{(R/a)^{7/2}}{\varepsilon_1 \text{watts/cm}^3}, \]

\[ \varepsilon_1 = \langle -T \frac{dT}{dr} \frac{q^3}{r^3} \rangle \]

\[ q \equiv \frac{r}{R} \frac{B}{B_{\text{pol}}(r)} \]

where \( Z \) and \( A \) are the effective ion charge and mass numbers, \( n_i \) is the ion density in cm\(^{-3}\), \( q \) is the "safety factor" equal to \( 2\pi \) divided by the rotational transform angle, \( T_i \) is the ion temperature in electron volts, \( B \) is the toroidal field in gauss, \( r \) and \( R \) are the minor and major plasma torus radii in cm. We have converted heat loss which is really a surface phenomenon to an equivalent loss per unit volume by calculating the flux of heat across the surface \( r = \) constant and dividing by the volume \( 2\pi R \cdot \pi a \). With a parabolic temperature profile, the use of \( r = a/2 \) will give \( \varepsilon_1 = 4 \sqrt{2}/3 \). \( T_i, n_i, \) and \( q \) should also be evaluated at \( r = a/2 \), but we will think of them as some sort of average values for ion temperature and density and for \( q \) in the plasma.

**Ohmic Heating.** In present tokamak experiments, the principal source of power input to the plasma is just the Ohmic heating from the plasma current. The question of anomalous resistivity has not yet been finally resolved, but it would appear possible to account for most of today's tokamak resistivity measurements by taking into account the neoclassical effects together with the enhancement of resistivity due to the presence of high-\( Z \) impurity ions.
The neoclassical effects on the resistivity appear for two reasons: first, the number of conduction electrons is reduced by the ratio of untrapped or conduction electrons to the total number of electrons, and second, the trapped electrons are stuck in the valleys of the magnetic field and they act, together with the ions, to impede the flow of the untrapped electrons. These two processes are about equally important, and together they modify the Spitzer conductivity by the factor $\varepsilon_2$ below.\footnote{1} Now the current density can be expressed in terms of the poloidal magnetic field,

$$j_n \approx \frac{B_{pol}}{2\pi R} = \frac{B}{2\pi Rq}$$

and the Ohmic heating power to the plasma can then be written

$$P_{\text{Ohm}} = 1.32 \times 10^{-2} \frac{Z_{\text{eff}}}{R} B^2 \frac{\log \Lambda}{q^2 T_e^{3/2}} \varepsilon_2 \varepsilon_3 \text{ watts/cm}^3$$

$$\varepsilon_2 \equiv \langle 1 - 1.95(\frac{r}{R})^{1/2} + 0.95(\frac{r}{R}) \rangle$$

$$\varepsilon_3 \equiv \langle 1 - \frac{J_d}{J_n} \rangle$$

$$\frac{j_d}{J_n} = \frac{4\pi n (kT_e + kT_i)}{B^2} q^2 (\frac{R}{r})^{3/2} \left[ 1.22 \left( \frac{\xi}{n} \frac{dn}{dr} \right) + \frac{0.35 r}{kT_e + kT_i} \frac{d(kT_e)}{dr} \right] - \frac{0.21}{kT_e + kT_i} \frac{d(kT_i)}{dr}$$
Here $Z_{\text{eff}}$ is the effective $Z$ of the plasma and $T_e$ is the electron temperature in electron volts. The reduction factor $\varepsilon_3$ represents still another neoclassical effect — a portion of the toroidal current in the tokamak plasma is due to a sort of diamagnetic current in the poloidal magnetic field, the famous "bootstrap" current ($j_b$). Since it is driven by pressure gradients rather than by the toroidal $E$ field, its contribution to $\eta j_o$ must be subtracted off. It should be mentioned that the numerical coefficients for both $\varepsilon_2$ and $\varepsilon_3$ are for a $Z = 1$ plasma.

We now have quantitative expressions for the two effects which are the dominant processes in laboratory tokamak experiments — power input by Ohmic heating and power loss by thermal conductivity. If we equate the expressions for these two powers, there occurs a most remarkable cancellation of terms leaving a very simple answer in terms of a ratio called $\beta_0$, defined

$$\beta_0 \equiv \frac{8\pi n_i (kT_i + ZkT_e)}{B_{\text{pol}}}$$

in which $B_{\text{pol}}$ is the poloidal magnetic field measured at the edge of the plasma, $r = a$. Now equating $P_Q = P_{\eta}$, we find

$$\beta_0 = 0.31 \frac{(1 + \alpha)}{\alpha^{3/4}} \left(\frac{Z R}{A a}\right)^{1/4} \left(\frac{\varepsilon_2 \varepsilon_3}{\varepsilon_1}\right)^{1/2}$$
where
\[ \varepsilon_3 \equiv \langle 1 - \frac{j_{\phi}}{j_{\phi}} \rangle \approx \langle 1 - 1.29 \beta_\theta \sqrt{\frac{\rho}{R}} \rangle, \]
\[ \alpha \equiv \frac{Z T_e}{T_i}. \]

In a representative evaluation, we use \( r = a/2 \), parabolic temperature and density profiles, \( T_e = 3 T_i \), \( Z = 2 \), \( A = 4 \), and \( R/a = 9 \) which are numbers typefying helium plasmas in Princeton's ST tokamak. Substituting in the above expression then gives the value \( \beta_\theta = 0.69 \). We can compare this with experimental values, some of which are plotted in Figure 3. Experiments show \( \beta_\theta \) lying in the range typically between 0.5 and 1, which is in remarkable agreement with the neoclassical calculation.

On the other hand, a more detailed look at the comparison of theory and experiment reveals some strong inconsistencies. First, as is well known, the neoclassical picture is unable to account for the observed electron heat loss. On the neoclassical basis, it is the ions which transport the heat to the plasma edge, just as we indicated earlier. Electron heat is to be delivered to the ion distribution via electron-ion collisions. But the classical rate for electron-ion thermalization is considerably longer than the observed energy confinement time. Additional channels for heat loss have therefore been proposed, such as the enhanced electron thermal conductivity coefficient in the "pseudoclassical" description.

Another discrepancy with theory occurs because the actual tokamak plasmas just referred to are still not quite in the full
neoclassical regime, but actually are in the "plateau" regime. For the same tokamak conditions used above to estimate $\beta_0$, the plateau thermal conductivity is only about half as large as the neoclassical thermal conductivity. This factor would increase the expected $\beta_0$ value by $\sqrt{2}$.

The indication then is that the actual experimental rate of heat loss is currently perhaps two to four times larger than the best theoretical value. In addition, we will shortly be talking about adding supplementary power to a tokamak plasma, and it is very possible that the plasma heating process will introduce its own anomaly in the heat transport rate. Therefore, in our power balance computations below, we will actually use a value for thermal conductivity which is equal to four times the neoclassical value.

Let me now describe rather briefly the remaining major heat loss and heat input mechanisms.

**Line Radiation.** The bulk of the energy radiated from current experimental plasmas comes from discrete spectral lines of partially-stripped impurity ions. However, this heat loss should become relatively much less important as we go to hotter cleaner plasmas, and we will not include it in our power balance calculations.

**Bremsstrahlung.** Bremsstrahlung radiation is emitted in the free-free transition as electrons are scattered by ions. The intensity in the bremsstrahlung spectrum falls off as $\exp(-hv/kT_e)$, therefore the radiation will lie mostly in the soft x-ray region and will not be reflected at the walls. The power loss per unit
volume due to bremsstrahlung may be written

\[ P_{\text{brems}} = 1.69 \times 10^{-32} z^3 \frac{n_i^2}{\text{eff}} T_e^{1/2} \text{ watts/cm}^3 \]

where \( T_e \) is in eV, \( n_i \) in cm\(^{-3}\).

**Synchrotron Radiation.** Like bremsstrahlung, synchrotron radiation also comes from a free-free electron transition. The electron acceleration in this case is due to the magnetic field, rather than to Coulomb scattering. The radiation is strongly peaked at the electron cyclotron frequency and its harmonics, and the transition rate is sufficiently great in a plasma that reabsorption is an important process. In fact, the usual approximation for the intensity of synchrotron radiation is to assume a black-body spectrum of radiation up to some cut-off frequency, typically expressed as a certain multiple of the cyclotron frequency. This cut-off appears where the plasma goes transparent, i.e., where the absorption length would be about equal to the radius of the plasma. Wall reflectivity may be expected to be quite high for this range of frequencies, however, so the effective absorption length is really the plasma dimension divided by the wall absorptivity. Moreover, the heat loss from the plasma is then most easily calculated as the heat absorbed by the walls, namely black-body or Rayleigh-Jeans radiation up to cut-off multiplied by the absorptivity. This calculation gives an effective power loss per unit volume from the plasma of

\[ P_{\text{syn}} = 1.64 \times 10^{-20} \alpha \frac{m^3 B^3 T_e}{R} \left( \frac{R}{a} \right) \text{ watts/cm}^3 \]
where $B$ is the toroidal field in gauss, $T_e$ is in eV, $R$ in cm, and $m$ is the harmonic number at cutoff, which we take equal to 6 (see Ref. 5). $\alpha$ is the absorptivity coefficient for the wall, multiplied by the wall surface area and divided by the plasma surface area. We will use $\alpha = 0.05$.

**Alpha-particle Heating.** Plasma heat to a working fusion reactor is mainly provided, of course, by the 3.52 MeV alpha particles which come from the D-T nuclear reaction. For the purpose of simple computations, we can make a reasonable fit to the averaged D-T cross sections $^6$ between $kT_i = 5$ keV and 20 keV using

$$
\langle \sigma v \rangle = 3.4 \times 10^{-57} T_i (-0.843 \ln T_i + 17.9) \text{ cm}^3/\text{sec},
$$

where $T_i$ is the ion temperature in eV. The heating power per unit volume supplied to the plasma can then be written

$$
P_\alpha = 4.8 \times 10^{-70} n_i^2 T_i (-0.843 \ln T_i + 17.9) \text{ watts/cm}^3.
$$

for the case of a 1:1 D-T mixture. $n_i$ is the total ion density in cm$^{-3}$.

**Supplementary Heating.** Finally, we will put into the power balance equation a term corresponding to the supplementary heat which we supply through radiofrequency heating, neutral injection, etc. For a toroidal plasma, we can write the power input per unit volume as

$$
P_{\text{sup}} = \frac{P_{\text{input}}}{2\pi R^3} \left( \frac{R}{a} \right)^2 \text{ watts/cm}^3
$$
where \( P_{\text{input}} \) is the total amount of supplementary heat in watts, and \( R \) is the major plasma radius in cm.

**Power Balance.** Having looked at the terms individually, we can now put them into the power balance equation, equating heat input to heat loss. We have simply

\[
P_{\eta} + P_{\alpha} + P_{\text{sup}} = P_{Q} + P_{\text{brems}} + P_{\text{syn}}
\]

Our expressions for each of these terms are given on the basis of equivalent heat per unit volume. We have glossed over all the subtleties of the radial dependencies, as well as the time variations. In preparing the following set of graphs, we have solved this zero-dimensional power balance equation for \( n_i \) as a function of \( T_i = T_e \), and as indicated above, with \( P_{Q} \) equal to four times its neoclassical value. Also, in our actual computations we have dropped the bootstrap term which, in the cases considered, would have decreased the asymptotic densities, but by less than 10%.

Before turning to the actual graphs, however, it is useful to write down the dependence of each of the terms in the power balance equation on the principal parameters of interest. If we extract from each of the terms its dependence on \( n_i, T, R, q, \) and \( B \), and represent the remaining part of the power term by a coefficient \( p \), we can write the power balance equation in the following form:
This relation will be helpful in the interpretation of the graphs, and also provides a general scaling law for steady-state power balance.

It is also illuminating to note that this power balance relation forms a simple linear equation in \( n_i^2 \). The condition \( n_i^2 \to 0 \) is noted on the graphs as the "synchrotron cutoff", and the condition for \( n_i^2 \to \infty \) is designated as "ignition".

**Dependence on Machine Size** \( (R) \). In the first two power balance plots, Figs. 4 and 5, we use machine size as a parameter, and draw the power balance curves for machines of various major radii, keeping the aspect ratio and other parameters fixed.

Heat loss by thermal conductivity, which is a diffusive process, is reduced relative to nuclear heating as we go to larger machines, and we see that curves for \( \text{Ze} = 1 \) corresponding to the largest machines reach ignition. If there were no thermal conductivity losses, ignition would have been reached at the "ideal ignition temperature" where bremsstrahlung is the only loss, just above \( kT_i = 4 \) keV. For the smaller machines, the thermal conduction loss is so great that the plasma never reaches ignition — at high temperatures the large synchrotron loss drives the asymptotic density to zero before the ignition temperature is reached.
With $Z_{\text{eff}} = 2$, the thermal conduction loss is increased by a factor of four, and the bremsstrahlung by a factor of 8. Under these conditions, Ohmic heating by itself is unable to bring even the largest device here to the ignition point.

**Effect of Supplementary Heating on a Laboratory Tokamak.** The next plot, Figure 6, shows the effect of different amounts of supplementary heating power added to the plasmas in a tokamak of major radius equal to 100 cm. We have reduced the magnetic field strength to 30 kilogauss, taken $q = 3$, and conservatively set $Z_{\text{eff}} = 2$ even in this deuterium plasma. This set of parameters would be an approximate description of current laboratory tokamaks. Operating on Ohmic heating alone ($P = 0$), the graphs indicate average temperatures around a kilovolt with densities around $10^{13} \text{ cm}^{-3}$. What is remarkable here, now, is the impressive change that can be achieved in the operating parameters with rather modest amounts of added power. With only 100 kilowatts of supplementary heating, the plasma temperature can reach 2 kilovolts at an ion density of $10^{13}$. With a megawatt of added power, temperatures greater than 5 keV should be available to us at the same densities. Although this achievement sounds very optimistic, let me recall our assumptions that the heat loss here is computed to be $2^3 = 8$ times the $Z = 1$ bremsstrahlung loss and $2^2 \times 4 = 16$ times the $Z = 1$ neoclassical rate for thermal conduction.

**Effect of Supplementary Heating on a Reactor Plasma.** Finally, we look at plots for a tokamak plasma with reactor parameters, Figures 7 and 8. While the $Z = 1$ plasma can be brought to ignition without supplementary heating, which we noted before, we
again see here that really impressive changes can be brought about with only relatively modest amounts of supplementary power. Ten megawatts is able to bring the $Z = 1$ plasma to ignition at almost full operating density, around $10^{14}\text{cm}^{-3}$. And even with the more conservative choice, $Z_{\text{eff}} = 2$, the addition of ten megawatts brings the reactor to ignition in following a trajectory in $n_{\text{i}}, T$ space that need never drop below a density of $3 \times 10^{13}\text{cm}^{-3}$. However, for a reactor of this size, which might produce 5,000 megawatts of electrical power, even 100 megawatts of supplementary heat would have to be considered a modest amount, and with 100 megawatts input, the critical density with $Z_{\text{eff}} = 2$ is moved up above $10^{14}\text{cm}^{-3}$.

In conclusion, we know from tokamak experiments that Ohmic heating has been very effective in bringing the plasma to temperatures in the one to two kilovolt range. Moreover, we have seen that it might be possible to bring a reactor plasma all the way to ignition just with Ohmic heating. However, what is of particular importance for this Symposium on Heating, is that really impressive changes in temperatures and densities achieved in the laboratory can be brought about with truly modest amounts of added heat, and similarly, that modest amounts of heat added to a reactor plasma make the problem of reaching ignition very much simpler.

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Fig. 1. Artist's sketch of Princeton reference design model fusion power reactor. Magnetic field at center line of plasma is 60 kG, R = 900 cm, a = 200 cm.

Particle trajectory in the \( r, \phi \) plane.

Fig. 2. Banana orbits in a tokamak. From B. B. Kadomtsev and O. P. Pugutse, in Reviews of Plasma Physics (Consultants Bureau, New York, 1970), Vol. 5, p. 249.
Fig. 3. Experimental data from the ST and T3-A tokamaks, showing approximate constancy of $\beta_\theta$. From reference 3.

Fig. 4. Power balance for tokamaks of various major radii. For $Z_{\text{eff}} = 1$. 

\[ \beta_\theta = \frac{n k T}{B_0 / \beta n} \]
Fig. 5. Power balance for tokamaks of various major radii. For $Z_{\text{eff}} = 2$.

Fig. 6. Effect of supplementary heating on laboratory-size tokamak.
Fig. 7. Effect of supplementary heating on reactor-size tokamak. For $Z_{\text{eff}} = 1$.

Fig. 8. Effect of supplementary heating on reactor-size tokamak. For $Z_{\text{eff}} = 2$. 
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