electron temperature is larger than the ion bulk temperature, and the effective drift velocity exceeds the ion sound velocity. Large electron density and temperature gradients in the piston region considerably increase the effective drift velocity $v_{d,\mathrm{eff}}$ for wave growth. Under these conditions the observed strong electron heating may be attributed to ion acoustic turbulence. Development of an ion tail and the approach to a marginal stability $v_{d,\mathrm{eff}} \approx v_{\mathrm{crit}}$ (Fig. 3) indicates stabilization of the ion acoustic instability by linear Landau damping. The ions are effectively heated by reflection from the magnetic piston. The degree of ion reflection increases with higher initial density.

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Transport during Turbulent Heating in a Tokamak[(a)]

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A model for the University of Texas turbulent torus is described. The prediction of the model shows a strong skin effect, and is in agreement with the experimental data if one saturates the Buneman instability by using the electron-trapping condition and the ion-sound instability by using the anomalous collision frequency $\nu^{A} \simeq 10^{-5}(T_e/T_i)(u/v_e)\omega_{pe}$, which *correlates* with ion tail formation. An interpretation of the laser-scattering temperature measurements is proposed.

Turbulent heating has been proposed as an alternative method to Ohmic heating because of hypothetically smaller losses which make the procedure appealing when considering schemes which aspire to reach ignition for a tokamak plasma. Previous experimental and theoretical work[1] does not emphasize on the one hand the formation of a skin current which limits the efficacy of the current-driven turbulence to heat the bulk plasma during the fast rise time of the current ($\tau <$ few μsec) for interesting densities ($n \gtrsim 5\times10^{12}/\mathrm{cm}^3$); on the other hand, neither does it emphasize the possibility of low-quality heating due to significant amounts of energy found in the tail of the particle distribution functions. The University of Texas turbulent torus,[1] through a more complete set of diagnostics, with a scaling similar to previous fast devices,[1] extends previous experimental research in the area of the current-penetration stage.

In this Letter, I present a theoretical model of the experiment during the fast rise time of the current, in which the turbulent transport is modeled after Liewer and Krall's work.[2] The numerical results for current-density profiles, Fig. 1(a), show the strong skin effect. The corre-

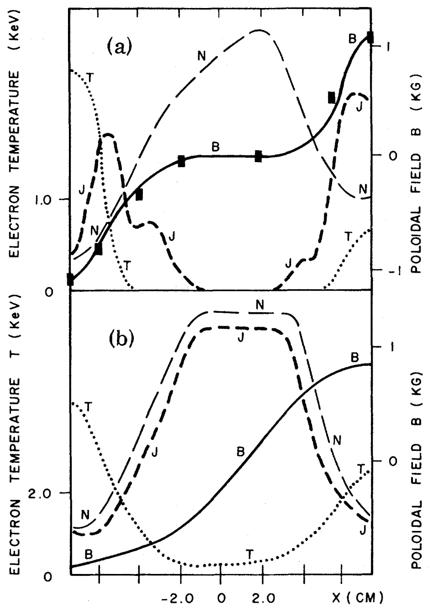

FIG. 1. (a) Line N, initial density profile for poloidal-field measurements i.e., with the presence of the magnetic probes during the preheat. The right-hand edge density is $4\times 10^{12}/\mathrm{cm}^3$; the peak value is $1.3\times 10^{11}/\mathrm{cm}^3$. Line B, poloidal-field profile at 1.2 μsec calculated with the value for ν^A from Mitchell *et al.* The squares indicate the experimental values. Line J, current-density profile at time 1.2 μsec; the peak value is 11.5×10^{12} statamperes/cm^2. Line T, temperature profile at 1.2 μsec. (b) Line N, density profile for laser temperature measurements, i.e., without the perturbation of the magentic probes, at time 770 nsec calculated with Caponi and Davidson's value for ν^A. The peak value of the density is $1.5\times 10^{13}/\mathrm{cm}^3$. The initial density (not shown) is slightly broader. Line B, poloidal-field profile at 700 nsec. Line J, current-density profile, at 700 nsec showing complete penetration. Line T, temperature profile at 770 nsec.

sponding poloidal magnetic field profile also shown in Fig. 1(a) is in good agreement with recent experimental values[3] when one saturates the ion-sound instability so as to agree with the simulation results of Mitchell *et al.*[4] Differing from previous work,[2] I note that the current penetration (or the poloidal-field profile) is fairly sensitive to the saturated value ν^A of the anomalous collision frequency for the ion-sound mode. Since different workers[4-7] have proposed different ν^A's, which can differ by two orders of magnitude between each other, it has been necessary to consider different versions in the present numerical analysis. However, for brevity, I report here only two of them, leaving the remaining cases for a future work.[8]

The theoretical model describes a plasma which supports a toroidal magnetic field and which is penetrated by a poloidal magnetic field, inducing a toroidal current.[8] The code self-consistently solves a system of equations for field and temperature profiles for given initial conditions (i.e., density and temperature of the species) and value of the poloidal magnetic field at the edge of the plasma as a function of time. This system of equations for modeling the transport during the fast rise time of the poloidal field, which can accommodate different regimes depending on the theoretical model used for ν^A, is

$$\frac{\partial n}{\partial t}+\frac{\partial n V_x}{\partial x}=0,$$

$$\frac{\partial n V_x}{\partial t}+\frac{\partial n V_x V_x}{\partial x}=\left((J_y B_z - J_z B_y)c^{-1}-\frac{\partial(\sum_j n_j T_j)}{\partial x}\right)(\textstyle\sum_j m_j)^{-1},$$

$$\frac{\partial J_y}{\partial t}+\frac{\partial V_x J_y}{\partial x}=\sum_j (E_y - V_x B_z c^{-1})\frac{e_j{}^2 n_j}{m_j}+\sum_j \frac{e_j R_{jy}}{m_j}, \qquad (1)$$

$$\frac{\partial J_z}{\partial t}+\frac{\partial V_x J_z}{\partial x}=\sum_j (E_z + V_x B_y c^{-1})\frac{e_j{}^2 n_j}{m_j}+\sum_j \frac{e_j R_{jz}}{m_j},$$

$$\frac{3}{2}\frac{\partial n_j T_j}{\partial t}+\frac{3}{2}\frac{\partial n_j T_{jx} V_{jx}}{\partial x}=R_{jT}-n_j T_j\frac{\partial V_x}{\partial x},$$

$$\frac{\partial B_z}{\partial x}=-\frac{4\pi}{c}J_y, \quad \frac{\partial E_y}{\partial x}=-\frac{1}{c}\frac{\partial B_z}{\partial t},$$

$$\frac{\partial B_y}{\partial x}=\frac{4\pi}{c}J_z, \quad \frac{\partial E_z}{\partial x}=\frac{1}{c}\frac{\partial B_y}{\partial t},$$

where $n=n_j$ is the density, e_j the electric charge, m_j the mass, and V_x the fluid velocity. The transport terms R_{jy}, R_{jz}, and R_{jT} will be described below. Since the frequencies characteristic of turbulent heating of interest in this experiment are much higher than the bounce frequency, we can ignore toroidal effects and use either cylindrical or slab geometry. Because the initial condition, i.e., the measured initial density profile, is not symmetric in the plane defined by the radial poloidal (X, Y) directions [see Fig. 1(a)], approximations must be made in the reduction of the asymmetric poloidal-field-profile data to compare with cylindrical-geometry calculated values. Thus I have chosen to use the slab geometry in this work for direct comparison with the experimental data, recognizing that in either case there are limitations to any procedure which is one-dimensional.

As the current penetrates, gradient-driven modes as well as the usual current-driven ones can become unstable, and anomalous penetration has been attributed to effects of the former.[9] We now note that the anomalous viscosity and anomalous thermal conductivity are the only transport coefficients that, in principle, can compete with the anomalous resistivity during the fast rise time of the poloidal field.[8] However, it can be shown that for the typical fast-time experiments it is sufficient to consider the effect of the anomalous resistivity due to the dominant modes[8] (Buneman and ion-sound modes). The physical picture[8] is that as these current-driven modes grow from thermal noise to saturation, particles interact with them and give rise to an anomalous friction force. For the calculations of the R's I will use quasilinear theory saturated to agree with theory and/or particle simulation. Thus, briefly, if $u > v_e$, then the Buneman parallel-flow mode is dominant and the contribution from this mode to the anomalous friction force is

$$-R_e^{\,B} = \nu^B n m_e \vec{u} = (2u/u^2) 2\gamma^B W^B,$$

where $\vec{u} \simeq \vec{J}/em$ is the drift velocity. Here γ^B is the growth rate and W^B is the wave energy in the fields, which satisfies the quasilinear equation

$$W^B(t) = (T_e/2\lambda_D^3)\exp\int^t 2\gamma^B(t')\,dt'.$$

The maximum value for the wave energy is estimated by using the electron-trapping[2] condition $W^B < \frac{1}{8} n m_e u^2$. With use of the conservation of energy for the electrostatic turbulence, it is straightforward to obtain the electron-heating-rate terms for this mode if ion heating is specified. If $v_e > u > c_s$, the ion-sound mode is dominant. Here, as noted before, two-dimensional computer simulations and theory indicate different values and descriptions for the saturated state; thus I needed to run different versions and report only two cases here. The first,[4] shown in Fig. 1(a), can be characterized as follows: In the linear regime we have

$$-R_e^{\,A} = m_e n u \nu_e^{\,A} = (9/c_s) 2\gamma^A W^A;$$

$$W^A = (T_e/4\lambda_D^3)\exp\int^t 2\gamma^A\,dt',$$

$$R_{Te} = (9u/c_s - 5) 2\gamma^A W^A.$$

Conservation of energy for the electrostatic turbulence leads to $R_{Ti} = 8\gamma^A W^A$. In the saturated regime we have

$$\nu_e^{\,A} = 10^{-5}(T_e/T_i)(u/v_e)\omega_{pe} \simeq (W^A/n_e T_e)\omega_{pe},$$

with

$$R_{Ti} \simeq n T_e \omega_{pe}(m_e/m_i)^{1/2}(u/v_e) W^A/n_e T_e,$$

and the electron heating rate is obtained by invoking conservation of energy by the electrostatic turbulence. The second case [5] [see Fig. 1(b)] is given for comparison and uses the condition $W^A \leq n T_e (T_i/T_e)^{3/2}$ for the saturated fields.

I turn now to the description of the numerical results and the comparison with the experimental data. The experimental arrangement is discussed elsewhere.[1] Previous experimental data[1] indicated a flat density profile for the initial plasma while new exhaustive density measurements[3] indicate a typical initial plasma characterized by a high-density region ($n > 10^{13}/\mathrm{cm}^3$) in some portion of the plasma relatively near the center surrounded by lower-density plasma. (The inclusion of the initial density profile makes necessary a revision of conclusions from all previous experimental work that assumed a flat density profile for similar plasma formation conditions and claimed total current penetration in the regime considered here.) Thus, the initial density profile for the computer run in Fig. 1(a) is the measured density profile in the vertical direction during the preheat (at 400 μsec) with the magnetic probes in the vertical direction. For the computer run of Fig. 1(b), I have chosen as the initial density the measured density profile without the perturbation of the magnetic probes, as is the case when one measures temperature profiles by laser scattering. The microwave-averaged density for both initial plasmas is $n \simeq 4.8 \times 10^{12}/\mathrm{cm}^3$. The effective edge of the plasma at ± 7.5 cm is due to the presence of a very low-density plasma in most of the annular region near the limiter radius ($X = \pm 9.5$ cm) which cannot be measured with the probe ($n \leq 3 \times 10^{11}/\mathrm{cm}^3$); this very-low-density region is confirmed by the poloidal-field measurements which show negligible current in this annular region. The measured magnetic fields at the edge, well represented as a function of time by $B_y(a,t) = B_n \sin \pi t/2\tau$ and $B_z(a,t) = 12$ kG, are the relevant boundary conditions for the fields used in the numerical studies. For both runs, Figs. 1(a) and 1(b), we have $\tau = 1.2$ μsec and $B_n = 1.1$ kG. The initial temperature profiles for electrons and ions are chosen to be flat with electron temperature equal to the measured value, 5 eV, and an extrapolated electron-ion temperature ratio equal to 3. (Essentially there are no free parameters in the code.) The dots in Fig. 1 indicate three-shot-average measured values of the poloidal field. As previously stated, it is clear that good agreement is found when we use the saturation value $\nu^A \simeq 10^{-5}(T_e/T_i)(u/v_e)\omega_{pe}$ of Mitchell *et al.*, as compared with use of that given by Caponi and Davidson, for example, shown in Fig. 1(b). Clearly, from the current-density profiles we see that in the first case the current does not penetrate into a region defined by a radius of 2 cm, while in the second case there is complete penetration, in disagreement with the measurements. Other values for ν^A, agreeing with other theories,[6,7] give more penetration than desirable to scale with the measurements. Another diagnostic is the 10° laser measurement of the temperature. I wish to show that values extrapolated from the temperature profile from Fig. 1(a) are consistent with the measurements. The description of the system is given elsewhere[10] and involves a scattered volume for the laser light from a region which samples a distance $\Delta X \simeq 2.5-3.0$ cm. The experimental data[3] are consistent with a bulk temperature for the electrons of 200 eV or less and a long, runaway tail in the direction of the drift with an effective tail temperature of about 5 keV, both measured at peak current (i.e., about 1.2 μsec) near the edge of the plasma ($X \simeq 6.5$ cm), both temperatures otherwise being very low. Recall that by definition the code temperature is the trace of the pressure tensor, i.e., $nT_e = \frac{1}{3}m_e \times \mathrm{Tr}\int f_e(\vec{V}-\vec{V}_e)(\vec{V}-\vec{V}_e)d^3V$. For the purpose of this Letter, I can simply fit the distribution function by the sum of two drifting Maxwellians, one representing the bulk with 75% of the electrons, and the other the tail with 25%. After invoking conservation of the proper moments, we obtain a bulk temperature $T_b \simeq T_e/6$ and a tail temperature $T_t \simeq 4T_e$, where T_e is the code value. When comparing with the experimental data one needs to note that the measured temperature is a weighted average over the scattered volume of the laser light, thus a spatial average over the distance ΔX. Considerations which are straightforward, but rather lengthy to be given here,[8] show that the code temperature T_e should be divided by a factor close to 1.5 when comparing with the spatial, laser-averaged experimental value, $\overline{T}_e$. This leads to $\overline{T}_b \sim T_e/9$ and $\overline{T}_t \sim 2.7T_e$. Thus, from the left-hand side of Fig. 1(a), we have near the edge $T_e \simeq 2$ keV; hence, $\overline{T}_b \simeq 200$ eV and $\overline{T}_t \sim 5$ keV. These values scale well with the measurements, which obviously is not the case when I use the data from Fig. 1(b).

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