transition at λ_1 being broadened and partially self-absorbed as shown in Fig. 2(a). The absorption of the outer layer is given in Fig. 2(b). The radiation emerging from such a plasma (core radiation passing through outer absorbing layer) is shown in Fig. 2(c). The anomalous features of a portion of the spectrogram obtained by Jaegle *et al.* for laser-produced aluminum plasma are also reproduced in Fig. 2(c) and are seen to be in good agreement with the predictions of the model. A slightly more complex plasma model, incorporating different absorptions for the continuum and line radiation to simulate a time-dependent outer-layer absorption [as suggested by the results of the laser-probe measurement shown in Fig. 1(b)], leads to an even better fit to the results of Jaegle *et al*.

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Plasma-Return-Current Heating by Relativistic Electron Beams with $\nu/\gamma \sim 10$*

D. Prono,† B. Ecker, N. Bergstrom, and J. Benford
Physics International Company, San Leandro, California 94577
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We report experimental results of hydrogen-plasma heating using very intense relativistic electron beams with ν/γ up to 10. The beam-plasma interaction process is found to be plasma-return-current heating with strongly anomalous plasma resistivity. Peak plasma energy is 1.5×10^{19} eV/cm^3, corresponding to $T_e = 2.2$ keV at $n_e = 6.7\times10^{15}$ cm^{-3}.

The use of intense relativistic electron beams to heat plasma has been investigated by several workers, primarily in order to apply the large beam energies available at high power to the goal of heating plasma to thermonuclear temperatures. Previous experiments[1,2] have generally utilized beam energies of $\lesssim 2$ kJ, with $\nu/\gamma \lesssim 2$. The measured plasma energy has been $\sim 10^{17}$ eV/cm^3. The interaction causing the heating has been ambiguous in some cases,[1] and identified as electron-electron beam-plasma interaction in others.[2] Theoretical studies indicate that with high-ν/γ beams, return-current heating will be the dominant plasma-heating process.[3]

We report here the results of initial experiments using a very intense beam (500 kA, 40 kJ, $\nu/\gamma = 10$). Plasma heating of 10^{19} eV/cm^3 is measured. We find that plasma energy gain and beam energy loss are due to return-current heating: (1) The measured total energy loss per transported beam electron is found to equal $e\int E_z\,dz$, where E_z is the independently measured, macroscopic, induced axial electric field which retards the beam; (2) the plasma energy determined from diamagnetic signals, W_m, agrees well with plasma energy density due to Ohmic heating by the plasma return current, independently determined from $W_{\rm Ohm} \equiv \int_0^{\tau_b} J_p(z,t)E_z(z,t)\,dt$, where J_p is the axial plasma current and τ_b the beam duration. (The contribution from azimuthal currents and fields is small.) Together these observations show that the beam-plasma interaction process is return-current heating.

The experiment consisted of injecting the intense relativistic electron beam through the transmission anode foil of the beam-generating diode, which was one end of a metallic hydrogen-filled cylindrical chamber of radius $r_w = 7.5$ cm. The neutral hydrogen fill pressure ranged from 30 to 1000 mTorr (atomic densities of 2×10^{15} to 6.7×10^{16} cm^{-3}, respectively). A 16-kG longitudinal magnetic field permeated both the 1.1 m chamber and the vacuum diode. No provision for plasma containment was made. The OWL II

Pulserad[4] provided a 40 kJ, 80- to 150-nsec pulse at peak currents of 200 to 500 kA, peak electron energy ~1 MeV, and with ν/γ ranging from 4 to 10. Beam radius was $a = 3.6$ cm. Local magnetic field probes located at $r_{pr} = 4.25$ cm measured $B_\theta(z,t)$ and $\Delta B_z(z,t)$ near the beam surface; B_θ gives the net current, $I_{net} = 2\pi r_{pr} B_\theta/\mu_0$. A Faraday cup and quartz stress gauge[5] (see below) located at the end of the chamber measured propagated beam current and beam electron energy, respectively, after propagation through the plasma.

Plasma current I_p is the difference of beam and net currents. The measured currents vary with distance z from the anode window: (1) Comparison of Faraday-cup and diode-current wave forms shows partial transmission of the injected beam (e.g., typically 40% at 100 mTorr H_2 and 75% at 1000 mTorr), and (2) comparison of B_θ wave forms from four different z positions shows I_{net} increasing downstream, implying radial current components, an effect which disappears as H_2 fill pressure is increased to 1000 mTorr. The beam current at each net-current (B_θ) probe position is needed for the I_p calculation. We use linear interpolation of the beam current based on the instantaneous values at the diode and Faraday cup. This is consistent with the approximately linear beam loss findings from earlier studies of high-ν/γ beams in neutral gas with applied B_z fields.[6] The beam-electron-loss mechanism remains unknown.

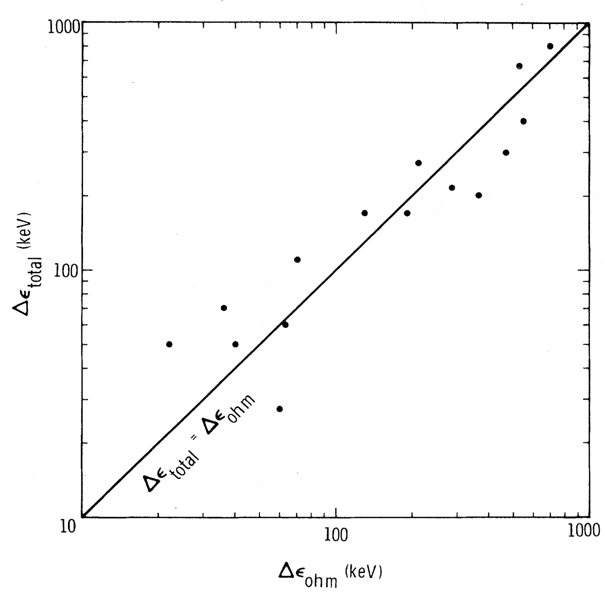

FIG. 1. Total beam-electron energy loss, $\Delta\epsilon_{total}$, versus beam-electron energy loss due to Ohmic-heating dynamics, $\Delta\epsilon_{Ohm}$.

Neglecting electrostatic fields and assuming radial uniformity of the net current density in $r \leq a$, the radial mean value of E_z in $r \leq a$ is $E_{z0} = L(\partial I_{net}/\partial t)$, where the system inductance is $L = (\mu_0/8\pi)[1 + 4 \ln r_w/a]$. We verify *a posteriori* that Ohm's law can be applied in simplest form; hence plasma conductivity is determined from $\sigma(t) = J_p(t)/E_{z0}(t)$, where $J_p = I_p/\pi a^2$. The heating rate is $\partial W_{Ohm}/\partial t = J_p(t)E_{z0}(t)$.

Our measurement of beam-electron energies involves a new application of the quartz stress gauge[5] which reverses the technique used in materials-response studies; here we use the piezoelectric stress signal and the known equation of state of the target plate to calculate the peak value of incident-beam-electron energy via Monte Carlo and hydrodynamic computer codes. Figure 1 shows the agreement between the total beam-electron energy loss measured using the quartz gauge, $\Delta\epsilon_{total}$, and the beam-electron energy loss to Ohmic heating, $\Delta\epsilon_{Ohm} \equiv e\int E_z dz$, where the integral is over the chamber length at the time of peak transmitted beam-electron energy. At all pressures the total energy loss is attributable to the retarding induced electric field, i.e., the Ohmic-heating process. Note that beam-plasma energy exchange via linear electron-electron two-stream instability would cause discrepancy between $\Delta\epsilon_{Ohm}$ and $\Delta\epsilon_{total}$ since macroscopic E_z is not generated by that process.

If there are no losses, the total energy present in the plasma as a result of return-current heating is W_{Ohm}. The horizontal axis in Fig. 2 shows W_{Ohm} computed from the measured $J_p(t)$ and $E_{z0}(t)$, versus neutral H_2 pressure via the vari-

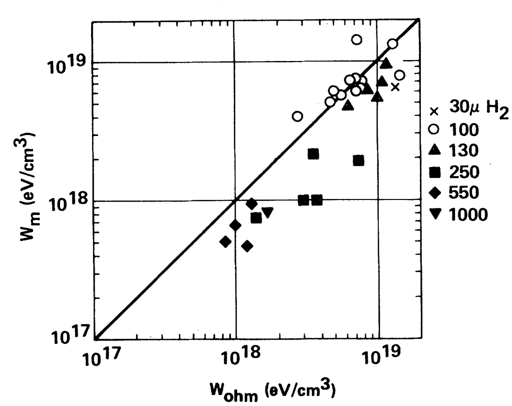

FIG. 2. Comparison of W_{Ohm} and W_m over the neutral H_2 pressure range studied.

a

b

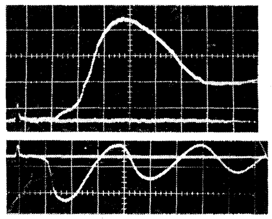

FIG. 3. (a) Net-current wave form, 100 mTorr H_2; 30 kA/div, 50 nsec/div. (b) ΔB_z wave form, 100 mTorr H_2; 1.0 kG/div, 100 nsec/div.

ous symbols. The decline of coupled energy as H_2 pressure increases from 100 to 1000 mTorr is the result of the higher conductivity attained at the higher pressures. [Ohmic heating after $t = \tau_b$ is relatively small because the driving field E_z has decreased. Also, as shown in Fig. 3(a) there is an abrupt decrease of $|\partial I_{\text{net}}/\partial t|$ at $t \lesssim 2\tau_b$ (when only plasma current remains) indicating a large increase in conductivity. This presumably results from the decay of turbulence.]

We use observed changes in axial magnetic field strength which accompany plasma heating [Fig. 3(b)] as an independent measure of the transverse energy density in the plasma, $W_{m\perp}$. At the higher pressures studied, the magnetic diffusion time, τ_m, is long compared with beam-pulse duration and therefore a frozen-field model is used in which the net pressure imbalance due to plasma heating causes radial expansion of the plasma column and the flux distribution within it. The expansion is quasistatically in pace with the increasing plasma energy because $\tau_b \gg a/V_A$, where V_A is the Alfvén velocity; ΔB_z stops growing within 15 nsec after τ_b, whereas inertia-dominated expansion would cause overshoot and continued ΔB_z growth long after the pulse. The static-fluid pressure-balance equation is therefore applicable; invoking flux conservation, we calculate $W_{m\perp}$ from ΔB_z measured outside the beam-plasma column. (Radial excursions are <1 cm, so that the beam-plasma-column boundary does not pass the probe.) At the lower pressures ($\lesssim$130 mTorr), lower conductivities during the beam pulse necessitate a diffusion model which also utilizes pressure balance but with no plasma expansion. Over the present parameter range the two models differ by at most 25% in $W_{m\perp}$ versus ΔB_z. For comparison with W_{Ohm}, $W_{m\perp}$ must be compensated for the longitudinal component $W_{m\parallel}$ (a factor of 1.5). Figure 2 compares the compensated value, W_m, with W_{Ohm}. At 100 mTorr, where heating is greatest, the agreement is quite good; at the other pressures $W_m \approx 0.6W_{\text{Ohm}}$, which is still rather good agreement considering the simplicity of the models used in these analyses.

Since W_{Ohm} is a time-integrated quantity while W_m is a nearly instantaneous measure of plasma energy density at the end of the beam pulse, the agreement implies little loss during the heating duration of 60 to 90 nsec. There is no provision for containment, and the low loss cannot be explained classically. We tentatively interpret it in terms of a high effective collision frequency ν^* impeding parallel thermal conduction and diffusion.[7] In the 100-mTorr case, for example, the measured conductivity σ (see below) implies that $\nu^* \approx \omega_{pi}$ and the plasma electron mean free path is of order 0.003 cm, whereas the classical value is of order 130 cm.

The decay of turbulence after the beam pulse (which we infer from rapid increase of σ) should, by these arguments, be accompanied by rapid plasma energy loss. This is in fact indicated by the observed fast decrease of ΔB_z [Fig. 3(b)]. The subsequent ΔB_z oscillation corresponds to radial bouncing.[8] The observed fill-pressure-dependent period agrees with theory if 100% ionization at late times is assumed, as is expected based on thermal ionization rate estimates.

Return-current-heating theory[3] predicts that $\pi a^2 W_{\text{Ohm}}/U_b = 2f - f^2$, where $f(t) = I_{\text{net}}(t)/I_{\text{beam}}(t)$ and $U_b = \frac{1}{2}LI_{\text{beam}}^2$ is the magnetic energy per unit length of the unneutralized beam. Evaluating just prior to τ_b, when f is well defined and heating is essentially complete, we find good agreement, $1 < \pi a^2 W_{\text{Ohm}}/[U_b(2f - f^2)] < 1.5$, over our entire range of data. Note that f varies by a factor of 14, going from 0.35 at 100 mTorr (at $t \lesssim \tau_b$) to 0.025 at 1000 mTorr.

Measured plasma conductivity rose during the beam pulse to typical values at $t = \tau_b$ of 350 mho/m at 100-mTorr H_2 fill and 7000 mho/m at 1000 mTorr, implying effective collision frequencies which are up to 10^5 times higher than can be accounted for by classical electron-ion or electron-neutral collisions.[9] This indicates a low-frequency turbulent wave spectrum. The time dependence of plasma density is unknown, but the broad range of density through which the plasma passes (10^{13} cm^{-3} to 10^{15} or 10^{16} cm^{-3}) leaves little doubt that the conditions for the current-driven electron-ion instability mode are satisfied during at least an initial portion of the pulse. However, computer calculations by Albright[10] with time-dependent T_e and T_i (but constant plasma density) indicate that ion-acoustic turbulence could account for the measured plasma energy only if ionization were $\lesssim 1\%$, which is very unlikely.[11] It is observed that the conductivity remains highly anomalous throughout the beam-pulse duration.

If the increasing plasma density stabilizes the electron-ion mode before the end of the pulse, low-frequency turbulence could be sustained by replenishment due to nonlinear coupling of the electron-electron streaming instability to ion density fluctuations.[12] In addition, the electron-electron mode could nonlinearly enhance low-frequency turbulence during the electron-ion unstable phase, increasing resistivity relative to the current-driven mode alone.[13] Young[14] has derived a nonlinear electron-electron–produced anomalous resistivity based on the oscillating two-stream instability, which is consistent with our resistivity data when 10%–20% ionization is assumed.

The agreement of W_{Ohm} and W_m over a wide range weighs strongly against the possibility that the high resistivity (and W_{Ohm}) results from a classical low-temperature ($\lesssim 1$ eV) bulk electron population while W_m results from a very energetic low-density population. Also, the endurance of a low-T_e bulk is inconsistent with Ohmic-heating rates of $\leqslant 60$ MW/cm^3. A lower-bound effective T_e is computed as W_{Ohm}/n_e assuming 100% ionization. This gives 2200 eV in the 100-mTorr case with highest W_{Ohm}.

In conclusion, we note that the appeal of the present interpretation is its consistency with our various independent and complementary measurements. Further experimental work is in progress to directly measure n_e, T_i, and T_e in order to delineate more conclusively the mechanism of the intense beam-plasma interaction.

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†Present address: Lawrence Livermore Laboratory, Livermore, Calif. 94550.

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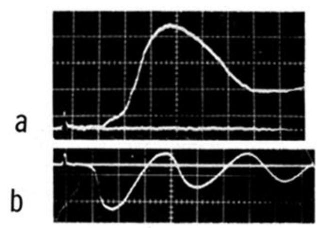

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