Plasma Heating by Intense, Relativistic Electron Beams

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A partially ionized helium plasma ($n_e = 5\times 10^{13}$ cm^{-3}) was heated by an intense ($\nu/\gamma \simeq 1$) relativistic electron beam (600 keV, 25–60 kA). Observations of the emitted line radiation and Thomson-scattered light indicate noncollisional electron heating rates consistent with electron-ion instability predictions, and total energy deposition consistent with electron-electron– instability saturation levels.

Advances in the design of electron-beam generators in the last few years have made it possible to produce electron beams delivering over 1.0 MJ of energy in times of about 100 nsec. Several possible uses for such high-power electron beams have been suggested.[1] Among the possibilities are intense radiation sources, including lasers that range from x-ray to microwave spectral ranges, gaseous-plasma heating and generation for controlled thermonuclear reactions and other purposes, and solid-target interactions for damage studies as well as controlled-thermonuclear-reaction–oriented programs similar to the laser-induced pellet-implosion scheme. A common requirement of most of these applications is an understanding of the interaction of the electron beam with a plasma medium.

The highest-power beams necessarily contain relativistic electrons (0.5 to 10 MeV). This means that the collisional interaction of the beam electrons with a low-Z background plasma is weak at plasma densities less than solid density. For example, in the experiment reported here the collisional-energy-loss mean free path for a single relativistic electron is greater than 10^8 cm. Consequently, the explanation of a strong interaction in a laboratory plasma requires a collective effect. Furthermore, because of the large energy densities involved, it will probably also require nonlinear effects.

Two broad categories of collective instabilities are likely to be important for intense relativistic electron beams in a plasma. First, we may have a two-stream electrostatic instability resulting from the beam electrons initially interacting with electron plasma waves in the background plasma.[2] A strength index for this instability with a cold beam is $S=\beta^2\gamma(n_b/2n_e)^{1/3}$, where the beam velocity is βc, $\gamma=(1-\beta^2)^{1/2}$, and n_b and n_e are the beam and plasma electron densities. The electron-electron instability is expected to transfer a maximum fraction of directed beam-energy density to plasma internal energy density for $S \gtrsim 0.1$.[3] Secondly, the induced field associated with the injected excess charge causes plasma electrons to stream back toward the accelerator and thus introduce the possibility of electron-ion streaming instabilities in the background-plasma return current.[4] Return-current heating is expected[5] to be most efficient for beams with current greater than the so-called critical Alfvén current $I_A = 17\beta\gamma$ kA, or as it is equivalently expressed $\nu/\gamma > 1$, where $\nu=(n_b\pi r^2)r_e$, with r_e the classical electron radius and r the beam radius.

Survey experiments[6-9] have been performed which show (1) beam energy loss from calorimeter or Faraday-cup measurements, and (2) a change in magnetic flux which is generally interpreted in terms of an increase of plasma thermal energy density. Some more direct measurements on plasma properties have also been made. For example, Altyntsev *et al.*,[6] using a neutral-atom detector, concluded that very little ion heating occurs for $\nu/\gamma < 1$, while Korn, Sandel, and Wharton,[8] using the same type of detector, observed energetic atoms from a plasma disturbed by a higher–current-density beam with $\nu/\gamma > 1$. We report here our measurements on the electron heating rate and how this relates to possible instability mechanisms.

The Triton accelerator contains a water-dielectric Blumlein transmission line which can be pulse charged to 1.3 MV by a Marx capacitor bank. The transmission line delivers maximum power to a matched load for 70 nsec. The experiments reported here are with a 500–600-kV charge and a diode which gives a peak current of 50–60 kA. Carbon cathodes of 5.0 cm diameter were used with 25-μm-thick titanium-foil anodes. Prior to firing the accelerator a plasma is created in a 5-m-long cylindrical glass vacuum chamber by z and θ current discharges. At the time the beam leaves the diode, z-directed magnetic field lines extend from the cathode through the entire plasma column. The magnetic field strength is 5 kG.

We have injected the beam into a partially ionized helium plasma ($n_e = 5\times10^{13}$ cm^{-3}, $p_0 = 100$ mTorr), and $\simeq 50$–70% of the beam energy arrived at the calorimeter 5 m from the anode. Beam transport efficiency appears to scale with density, similarly to earlier survey work. Streak photographs show the beam to be well centered in the discharge tube and 5 cm in radial extent. Coincident with the expected arrival time at 3.72 m a low level of microwave emission ($\lambda \leq 3$ cm) is detected. After 20 nsec a large narrow spike of microwave emission is observed followed by emission lines of He I ($\lambda = 4922$ Å) and He II ($\lambda = 4686$ Å). The microwave emission has a narrow bandwidth approximately centered on the relativistic electron cyclotron frequency. Thomson-scattering measurements indicated that the electron density increased to about $(7\pm3)\times10^{14}$ cm^{-3} within 100 nsec after the beam transit and stayed at that value for several hundred nanoseconds thereafter. At these late times the electron temperature was about 5 eV. At earlier times and during the beam passage the scattered light signal was too small with respect to the background radiation level for quantitative measurements, but qualitatively it indicated higher temperatures and lower electron densities.

At these temperatures and densities the plasma electrons are severely affected by ionization and radiation losses. In order to analyze our observations and to determine the energy transfer rate to the plasma electrons, the experimental observations were compared with a model described by the rate equations for ionization and excitation of helium atoms and ions along with the necessary energy equation. We assume that the rate of change of the internal energy density of the plasma ($\frac{3}{2}n_eKT_e$) is given by

$$\frac{\partial}{\partial t}(\tfrac{3}{2}n_eKT_e) = \dot{Q} - \sum_{j=0}^{1} E_j \frac{\partial n_j}{\partial t} - \sum_{j=0}^{1} E_{j,2}{}^* \frac{\partial n_{j,2}{}^*}{\partial t},$$

where $\dot{Q}$ represents the power-density input to the plasma electrons from the beam, E_j and n_j are the respective ionization potential and density of the jth ionization stage, and $E_{j,2}{}^*$ and $n_{j,2}{}^*$ are the first excitation potential and level-population density for the jth ionization stage. The use of a temperature here assumes thermalization on a time scale shorter than the ionization time. The ionization and excitation rate equations are

$$\partial n_{j+1}/\partial t = n_e n_j \langle \sigma_j V\rangle,$$

$$\partial n_{j,n}{}^*/\partial t = n_e n_j \langle \sigma_{j,n}{}^* V\rangle - A_{j,n} n_{j,n}{}^*,$$

$$n_e = n_1 + 2n_2,$$

where $A_{j,n}$ is the transition probability from the nth level. Because of the short time scales under consideration here we have neglected dynamical plasma motion, heat conduction, recombination, trapping of resonance line radiation, and deexcitation of highly excited states. The ionization cross sections σ and the rate coefficients $\langle\sigma V\rangle$ averaged over Maxwellian distribution functions were obtained from Lotz.[10] The excitation rates[11] used average Gaunt factors taken from Van Regemorter.[12] In addition to computing excitation to the $n=2$ levels to account for resonance-line radiation losses, we also computed the population densities for $n=4$ assuming excitation from the ground state in order to compare the model predictions with experiment for the time dependence of the 4922-Å He I and 4686-Å He II line radiation. It should be noted that the rapid ionization cannot be accounted for by electron impact from beam electrons because of the low reaction rate. Nor can the results be explained by the drift energy of the return-current electrons, because their energy is too low, $E_d \approx 0.5(n_b/n_e)^2 \times m_0c^2 < 9$ eV.

We have varied Q, the power-density input, in order to compare the model and our observations. Several possibilities exist for the power-density deposition term $\dot{Q}$. Return-current heating may be estimated in terms of a resistivity, $\dot{Q} = j^2\eta$, where j is the current density and the resistivity $\eta = (m_e/n_e e^2)\nu$. The collision frequency ν may be divided into electron-ion and electron-atom terms, and an effective collision frequency due to instabilities. In computing this heating rate we have assumed complete current neutralization and used the beam-current wave form at the observation point in the plasma, which has a peak value about 50% of the diode current. In Fig. 1, we have drawn curves showing the time dependence of the plasma electron density using several choices for ν. It can be seen that elastic collisions alone are not sufficient to explain the observed density increase. In addition, the time dependence of the line-radiation output is not as observed. Observed and calculated time histories of the He II 4686-Å line are shown in Fig. 2. Current neutralization requires that the plasma drift velocity $v_d \cong (n_b/n_e)\beta c$. For all but the first

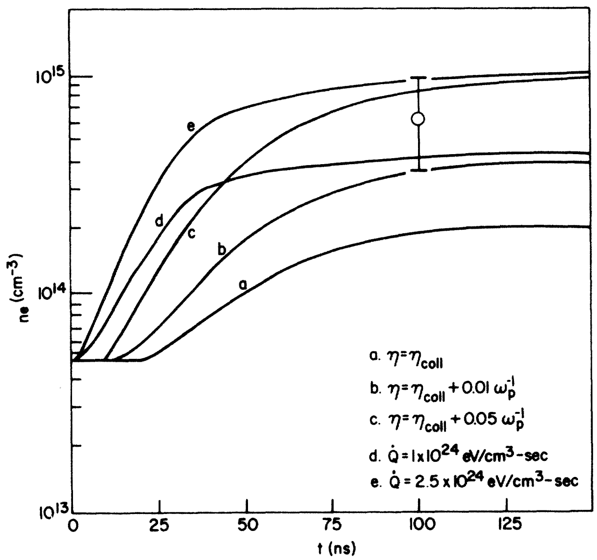

FIG. 1. Calculated electron-density variation as a function of time for different possible electron heating rates. Curve a, for Joule heating with a resistivity based on equilibrium plasma electron-ion and electron-atom collisions. Curves b and c include additional resistivity due to ion acoustic turbulence. Curves d and e, for a constant power-density input to the plasma electrons which lasts for 15 nsec. The experimentally observed range in the resulting electron density due to shot-to-shot variations is indicated by the bars above and below the circle. Similar results are obtained for later times up to 300 nsec.

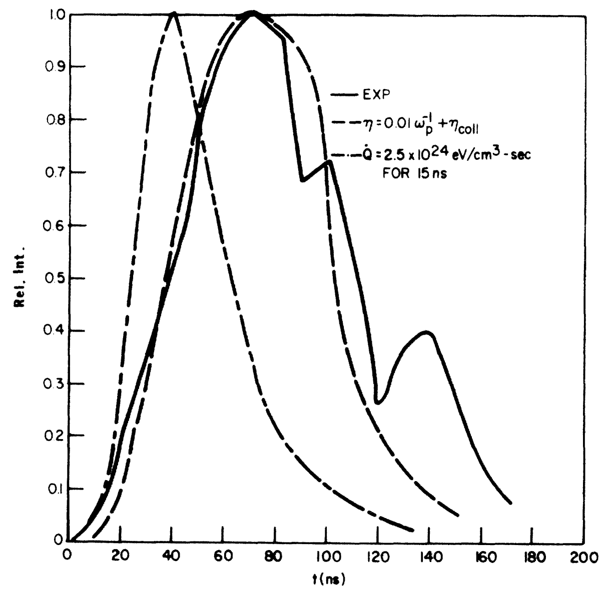

FIG. 2. Computed and observed time histories of the He II 4686-Å line intensity. Curves show agreement between the experiment and calculations using Joule heating based on a turbulent ion acoustic resistivity. Less agreement is seen for rapid energy deposition for 15 nsec, corresponding to the emission of X-band microwaves.

few nanoseconds the drift speed is greater than the ion acoustic speed and less than the electron thermal speed. This suggests using a resistivity based on ion acoustic turbulence. Various estimates[13,14] of the resistivity have been made and usually the resistivity falls within the limits $10^{-2}\omega_p^{-1} \lesssim \eta \lesssim 10^{-1}\omega_p^{-1}$, where ω_p is the electron plasma frequency. Inspection of Figs. 1 and 2 show that there is agreement between the model with $\eta = \eta_{coll} + 0.01\omega_p^{-1}$ and the experimental results for both the energy deposition and the variation of the ion line radiation with time. Substantially similar results are obtained by assuming a constant rate of energy deposition equal to 5×10^{23} eV/cm^3 sec which lasts for the length of time that the beam is on (90 nsec). This gives an energy of 50–70 J dissipated in the plasma electrons which results in a peak electron temperature of about 100 eV with the bound-state losses or about 750 eV with a fully ionized hydrogen plasma of the same initial density.

One might also assume that the burst of microwave radiation is indicative of the time history of energy deposition. In Figs. 1 and 2 we have also plotted curves for a constant $\dot{Q} = 2.5 \times 10^{24}$ eV/cm^3 sec lasting for 15 nsec. Although the total energy deposition (as evidenced by the increase in electron density) appears correct, the time variation of the He II 4686-Å line is not right.

The initial strength index for the electron-electron instability is $S \cong 0.17$. Using Thode and Sudan's estimate[3] of the fractional energy transfer, $W = \frac{1}{2}S(1+S)^{-3/2}$, we would expect about 7% of the beam energy density to appear in turbulence and an appreciable fraction of this to be thermalized. Comparison with theoretical predictions for the heating rate requires more detailed measurements of the beam temperature because the linear growth rates are temperature dependent. Our results thus appear consistent with either type of instability or perhaps both occurring in the plasma. Positive identification of unstable modes awaits more detailed measurements.

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High-Temperature Resistance Minima in Concentrated Pd-Ag and Ni-Cu Alloys

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The temperature dependence of the electrical resistivity of $Pd_{60}Ag_{40}$ between 4.2 and 900 K is shown to contain the same anomalies as previously found in Ni-Cu alloys in the midrange of composition. For Pd-Ag these effects cannot be due to scattering at giant polarization clouds, and are attributed to a temperature-dependent decrease of the impurity resistivity due to a reduction in *s-d* scattering with increasing temperature.

Three years ago Houghton, Sarachik, and Kouvel[1] reported that, at room temperature and above, the electrical resistivity ρ of concentrated paramagnetic Ni-Cu alloys shows a temperature dependence that is most unusual. Above ~200 K, in all alloys containing between about 50 and 62% copper, ρ *falls* with increasing T with the formation of a set of resistance minima at ~600 K in each case. This anomalous decrease in ρ they associated with the presence of giant polarization clouds that had previously been shown to exist in such alloys by neutron-diffraction measurements.[2] The suggestion was that a substantial fraction of ρ at low temperature can be attributed to electron scattering from these giant moments, but that as T rose and the clouds "dispersed" this form of scattering fell away. At the same time this reduction in ρ was overtaken by an increasing phonon resistivity, and this, qualitatively, was the origin of the resistance minima. Recently, Levin and Mills[3] have reported a detailed calculation of ρ based on a model of giant spin clusters, and discuss the internal dynamics of the clusters necessary to lead to the experimental results.

However, some further recent measurements by Houghton, Sarachik, and Kouvel[4] have cast doubts on this interpretation. They find that in Ni-Rh alloys close to the critical composition—a series in which there is again ample evidence of giant polarization clouds—there is no suggestion of a decrease in ρ at any temperature between 2 and 700 K.

We should now like to report the existence of a resistance minimum at ~700 K in the alloy $Pd_{60}Ag_{40}$ – that is, in an alloy that, while chemically equivalent to Ni-Cu, has magnetic properties that are quite different. Giant polarization clouds can only form in a metal that is *almost ferromagnetic*, and this is certainly not the case in Pd diluted by 40% Ag. Our data are shown in Fig. 1 together with measurements on a $Ni_{35}Cu_{65}$ sample (i.e., an alloy of similar resistivity that is paramagnetic at all temperatures) for pur-