version is left in the medium on the $4P$-$3D$ transition. This explains that the system can now superradiate after a longer delay at 9.10 μm.

The superradiant effects described in this Letter can have interesting applications. They should be quite generally observed in Rydberg levels which are very close to each other and should easily superradiate on infrared or microwave transitions having a small Doppler effect and hence a very low threshold for superradiance. Extension of our experiments to the visible is also very promising. The major drawback in this case would be the large Doppler effect at short wavelengths. This effect could be reduced by narrowing the width of the pumping dye lasers in order to select a narrow velocity group in the Doppler profile of the pumping transition.

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Observation of Collisionless Heating and Thermalization of Ions in a Theta Pinch*

R. J. Commisso† and Hans R. Griem
Department of Physics and Astronomy, University of Maryland, College Park, Maryland 20742
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Spectral line profiles from impurity ions in a small theta pinch have been measured and are interpreted as ion-velocity distribution functions. For antiparallel driving and bias fields, the ions attain a Maxwellian distribution corresponding to a perpendicular temperature of 3–4 keV in a time short compared to particle self-collision times. Thermalization is less rapid for parallel magnetic fields.

Achievement of deuterium and tritium ion temperatures over ~ 5 keV is critical for thermonuclear fusion.[1] In a theta-pinch, imploding electrons and ions acquire the same velocity, because of charge-separation effects.[2] This provides an efficient mechanism for putting directed kinetic energy into ions. In addition, for fusion reactions of the thermonuclear type to occur, the directed ion energy must be randomized.

We measured Doppler-broadened profiles of spectral lines from various oxygen and carbon ions which are natural contaminants and which radiate at different times due to transient ionization. These time-resolved profiles may be interpreted as the (unnormalized) ion velocity distributions. The experiment was performed on a small, 17-cm-diam, 15-kJ theta-pinch device,[3] with peak driving magnetic field $B_z \simeq 17$ kG and quarter period 2.25 μsec. The pinch produced a plasma from an 11-mTorr filling pressure of H_2, and the plasma was diagnosed using magnetic probes and Thomson scattering.[2] Peak electron (proton) temperatures and densities were ~ 250 eV (~ 900 eV) and ~ 6×10^{15} cm^{-3}, respectively.

A 0.5-m monochromator was used to scan the spectral line profiles on a shot-to-shot basis. The instrument function was nearly Gaussian with an instrumental width of 0.53 Å, always much less than the measured widths. Measurements were made both side-on (perpendicular to B_z) and end-on (parallel to B_z). Side-on, the monochromator viewed a 1-cm-high region of the midplane of the coil, centered on the coil axis; end-on, a 3-cm region. In most cases the total line intensity was monitored independently and used to normalize the data. The profiles were obtained for times between 0.61–1.5 μsec after the main bank discharged, i.e., after the implo-

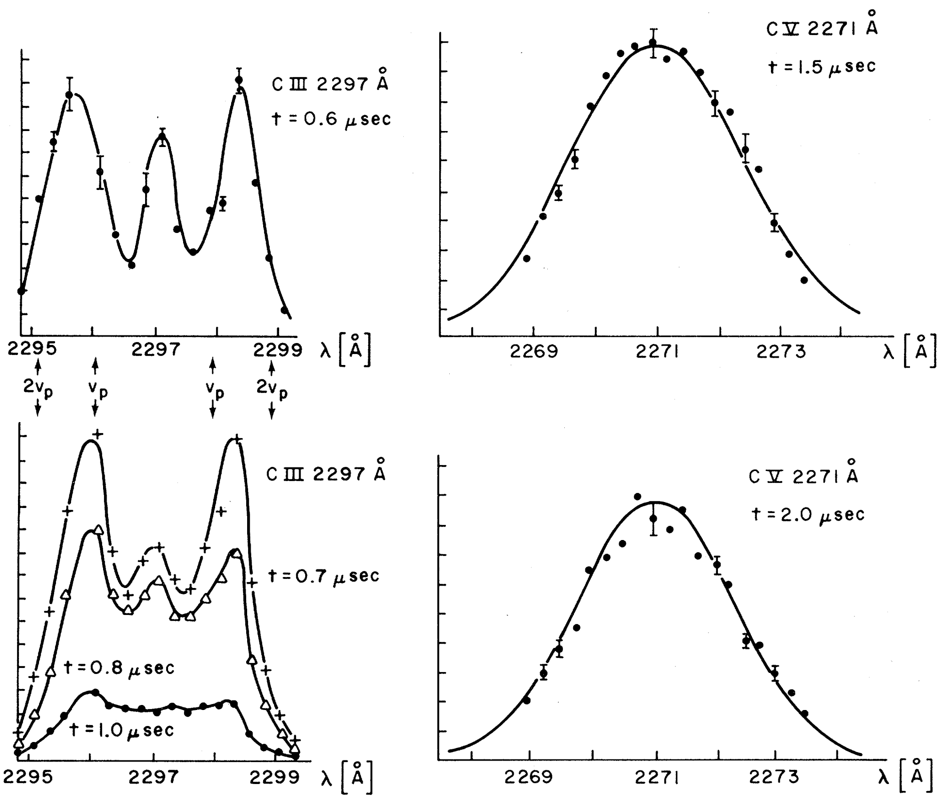

FIG. 1. Relative intensities of the C III 2297 Å and the C V 2271 Å lines at various times during the discharge, plotted against wavelength, for antiparallel fields.

sion phase ($\lesssim 0.5$ μsec).

Figure 1 shows the C III 2297-Å and the C V 2271-Å line side-on profiles. The C III 2297-Å line exhibits structure in the form of peaks (or shoulders) located symmetrically on either side of line center. These peaks correspond to Doppler-shifted radiation from C III ions moving at a speed v, $v_p \lesssim v \lesssim 2v_p$, where v_p is the measured[2] piston speed. The profile evolves in time from a triplet to a flat-topped structure. The error bars indicate the standard error of the mean. The triplet structure filling in indicates that the directed kinetic energy from the implosion is being randomized.

The profile of C V 2271 Å at later times is completely different. A Gaussian profile fits the data well. Thus, the C V ions have attained a velocity distribution which is very close to Maxwellian. The half-width at 1.5 μsec corresponds to a perpendicular temperature of $\simeq 4.3$ keV. At 2.0 μsec, the temperature has decreased to $\simeq 3.4$ keV.

This measurement was also done for oxygen ions, using the 3736.8-Å line of O IV and the 3811.4-Å line of O VI. The triplet structure is evident with the O IV profiles, consistent with the C III data. At 1.2 μsec, both the O IV profile and the O VI profile are flat-topped. At 1.5 μsec, the O VI profile is Gaussian again, at a temperature of $\simeq 3.1$ keV. The temperature is approximately the same at 2.0 μsec.

End-on profiles of C V 2271 Å were obtained at 1.5 and 2.0 μsec. They fit a Gaussian fairly well, with half-widths corresponding to parallel temperatures of $\simeq 2.2$ keV at both times. Care must be taken in interpreting this data, since the light intensity is integrated over the line of sight. If temperature gradients exist at both ends of the plasma, the measured profile is actually a superposition of Gaussians. This makes the observed width smaller, thereby giving an underestimate of the temperature. The problem is not as serious for the side-on measurements, since the ion gyroradii are comparable to the plasma radius.

Figure 1 indicates that (impurity) ions are being both swept up and, to varying degrees, reflected from the imploding piston. Then the ordered energy of directed motion is converted into random energy of thermal motion, the thermalization being completed ~1 μsec after the implosion. Assuming that all the directed kinetic energy acquired by the ions during the implosion is converted into thermal energy, one can estimate bounds for the final ion temperature. Ions acquire directed energy corresponding to speeds between v_p and $2v_p$, giving

$$\tfrac{1}{2}m_i v_p^2 \lesssim kT_i \lesssim \tfrac{1}{2}m_i(2v_p)^2. \qquad (1)$$

Here, m_i is the ion mass and it has been assumed that the ion velocities are two dimensional. For carbon, this yields 1 keV$\lesssim kT_i \lesssim$4 keV, to be compared to the measured temperature of $\simeq$3–4 keV. Actually, only a very small fraction of the ions have speed $2v_p$, making the upper bound rather extreme, and so the mechanism responsible for the rapid thermalization of the ions must also heat them (i.e., increase their mean kinetic energy). A similar estimate can be carried out for O VI, with essentially the same results.

Spitzer[4] has estimated the time scale, τ_c, for an arbitrary velocity distribution function of a group of particles to approach a Maxwellian distribution by virtue of the particles colliding with each other. For ions, τ_c is the characteristic collisional time scale in this experiment. It has a value for the measured C V and O VI temperatures, assuming $n_{\mathrm{C\,V,O\,VI}} \simeq 0.02 n_e$,[5] of $\tau_c \gtrsim 10$ μsec. This is much longer than the actual time of thermalization. Thus, some collisionless mechanism must govern the observed time evolution of the profiles. The additional heating of the ions is also faster than classical collision times.

The C V 2271 Å was also scanned with a *parallel* bias field, side-on. Figure 2 shows this profile at 1.5 and 2.0 μsec. It is not Gaussian, in contrast to the corresponding profiles in the antiparallel case. This indicates a less rapid thermalization in the parallel-bias-field case. The mechanisms responsible for the rapid thermalization and heating observed with antiparallel fields, therefore, appear to be absent or much weaker with parallel bias field. This may in part be due to the presence of a neutral layer (region of zero magnetic field) in the antiparallel case. While preferential heating with reverse bias field has long been known,[6] the results presented here also give evidence of more rapid thermalization.

The average impurity-ion gyroradii are large enough for the ions to bounce back and forth across the plasma column.[7] This could provide the randomization mechanism if the bounce time

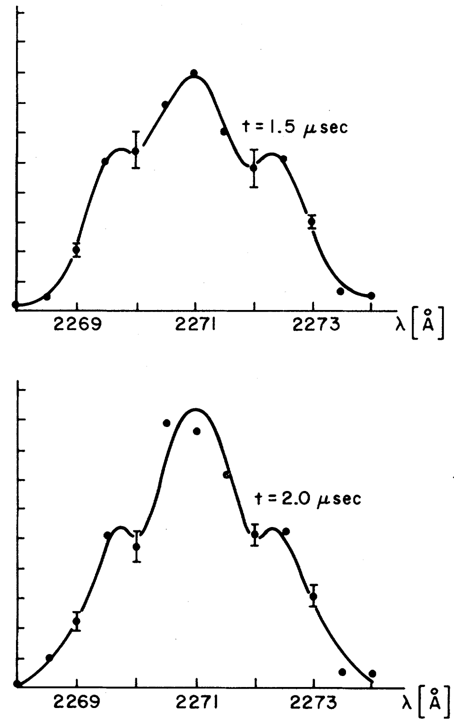

FIG. 2. Relative intensity of the C V 2271-Å line plotted against wavelength, for parallel fields.

was sufficiently small. The thermal speed of oxygen and carbon at temperatures of 3–4 keV is $\approx 2-3\times 10^7$ cm/sec, allowing ~5 bounces in ~1 μsec, which is probably sufficient to provide the observed thermalization.

However, this mechanism neither explains the observed additional heating of the impurity ions, nor the rapid thermalization observed end-on. Thus, a different collisionless mechanism suggests itself—that of wave-particle interaction. That is, unstable collective oscillations, deriving their energy from some source of free energy in the plasma, act to rapidly alter the velocity distribution functions of the ions. Besides instabilities peculiar to the neutral layer, three such instabilities, whose estimated growth rates are large enough and whose parameter regimes are appropriate for this experiment, may be considered.

One source of free energy to drive instabilities is the counterstreaming of ions across the plasma column. Papadopoulos *et al.*[8] have considered ion heating due to an ion-ion two-stream instability perpendicular to a magnetic (compressed bias) field. This instability essentially transforms the drift energy into thermal energy. Instabilities may also be fed by the free energy available from the relative drift of electrons with respect to ions in the sheath region.[2] Davidson and Gladd[9] have examined the anomalous resistivity and heating associated with the lower-hybrid-drift instability in the regime where $v_E \lesssim v_{\mathrm{Th}i}$, where $v_{\mathrm{Th}i}$ is the ion thermal speed and v_E is the azimuthal component of the $\vec{E}\times\vec{B}$ drift velocity. These waves propagate perpendicular to the magnetic field (i.e., azimuthally). Finally, the observed anisotropy in the C V temperature may also be a source of free energy. Davidson and Ogden[10] have studied transverse electromagnetic perturbations, propagating parallel to a magnetic field, which can become unstable in the presence of an ion-temperature anisotropy, $T_{i\perp}/T_{i\parallel}>1$. Numerical results[10] for this electromagnetic ion-cyclotron instability show that after ~10 maximum growth times (~1 μsec) the ratio $T_{i\perp}/T_{i\parallel}$ seems to saturate at ~2.

The data and analysis presented here show that collisionless mechanisms play a major role in the ion heating and thermalization in theta-pinches with densities two orders of magnitude higher than in earlier work on turbulent heating of hydrogen and argon ions.[11] One might add that impurity ion temperatures substantially larger than deuterium ion temperatures have also been observed in a dense plasma focus.[12]

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†Present address: Los Alamos Scientific Laboratory, Los Alamos, N. Mex. 87545.

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Plasma Heating by the Dissipative Trapped-Electron Drift Instability*

Nicholas A. Krall
Laboratory for Applied Plasma Studies, Science Applications, Inc., La Jolla, California 92037

and

Paulett C. Liewer
Department of Physics and Astronomy, University of Maryland, College Park, Maryland 20742
(Received 17 December 1975)

It is shown that energy loss in tokamaks due to cross field diffusion caused by the dissipative trapped-electron drift instability is compensated in part by the heating associated with this mode. Thus the energy loss rate presently predicted for tokamaks in the parameter range of this mode may be substantially overestimated.

It is generally accepted that instabilities driven by the diamagnetic drifts of trapped electrons and ions in tokamaks[1] can, in principle, strongly affect particle diffusion and heat conductivity in the relatively low-collision-frequency limit appropriate to controlled-fusion reactors. Indeed a great many of the present tokamak scaling predictions are based on a diffusion coefficient $D_\perp \cong \gamma/k^2$, where γ and k are the growth rate and wave number of trapped-particle drift instabilities.[2] However, diffusion is only one aspect of transport, and an instability which produces significant diffusion may be expected to contribute also to plasma heating and to the electrical conductivity. If the heating due to the mode is neglected, the energy loss predicted due to diffu-