Observation of Fast Neutrals Projected from a Coaxial Gun

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Xenon atom and ion densities were measured downstream from a coaxial plasma gun using vacuum ultraviolet emission, absorption spectroscopy, and a Langmuir probe. A pulse of fast neutrals was observed to emerge, in addition to a faster pulse of ions.

A spectroscopic study of the plasma projected from a coaxial plasma gun is described in this Note. This study has resulted in direct evidence of accelerated xenon neutrals in the plasma beam. Simultaneous observation of the ionized species has served to correlate these two species and has provided insight into the origins of both.

A Jarrell–Ash 1-m normal incidence concave grating vacuum spectrometer was used both as a scanning monochromator and a photographic spectrometer. The optical arrangement is shown in Fig. 1. The background light source used in these studies consisted of a water-cooled quartz capillary through which xenon gas flowed. Excitation was achieved in a slotted cavity driven by a 2.146 kMc/sec, 800 *W* oscillator. In addition, a spherical Langmiur probe of 0.75 mm diam. was used in order to estimate ion number densities. The spherical probe was prepared by melting thin copper wire in a flame, as described in detail elsewhere.[1]

Details of the plasma gun are also presented elsewhere,[2] but for the purposes of this discussion the xenon was fed to the discharge in discrete pulses in such a way that a major fraction of the pulse was in the interelectrode region at the time of the discharge. The discharge reached peak powers the order of 10^7 W and persisted for the order of 10^{-5} sec. Mass loading was such as to permit average particle velocities the order of 2–5 $\times$ 10^4 m/sec with the available energy per pulse. It should also be noted that the gun was not operating in its optimum mode during the course of these measurements.

A line absorption technique utilizing the XeI line at 1470 Å was applied to these studies. The basic relationship used is

$$I(\nu)/I_0(\nu) = \exp[-k(\nu)nL], \tag{1}$$

where $I(\nu)/I_0(\nu)$ is the transmissivity at a frequency ν, k is the absorption coefficient, n is the number density of absorbers, and L is the absorption path length. The exact expression for k is complicated, but under the assumption that the emission line is much narrower than the absorption line, a good approximation to the integrated transmissivity is given by

$$I/I_0 = \exp(-k_0 nL), \tag{2}$$

where I/I_0 is the transmissivity observed in practice with spectrometer slits set for a spectral range much greater than the width of the emission line from the background illuminator, and

$$k_0 = \pi^{1/2} r_e c f \delta_{DA}^{-1},$$

where r_e is the classical radius of the electron, f is the oscillator strength of the transition, and δ_{DA} is the Doppler width of the absorber.

An approximate velocity distribution function may be constructed from these data, in analogy with a technique developed for a multigridded particle probe.[3] The expression used for these purposes is

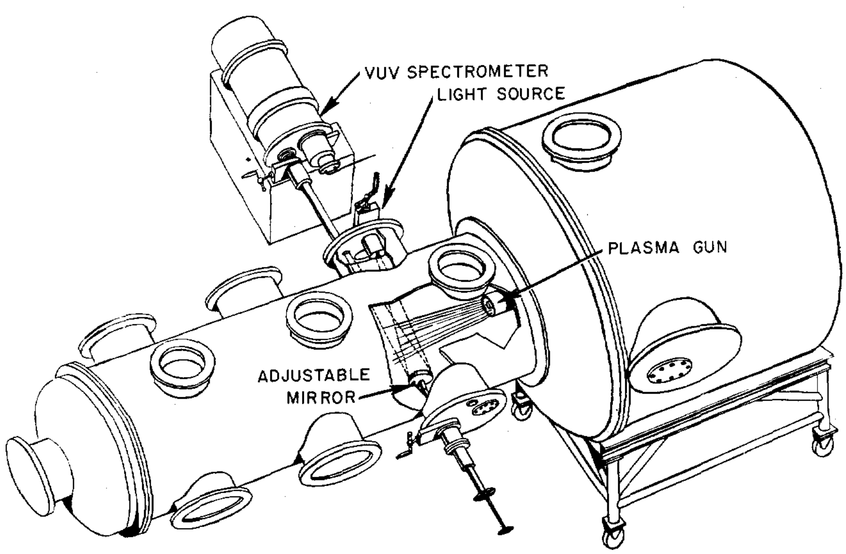

FIG. 1. Optical arrangement for vacuum ultraviolet measurements.

$$\frac{dN(v)}{dv} = -\left(\frac{dN(t)}{dt}\right)\left(\frac{dt}{dv}\right) = \left(\frac{nAD}{t}\right)\left(\frac{t^2}{D}\right) = ntA$$
$$= -(\pi^{1/2}L/16r_ecf)\,\delta_{DA}t\ln(I/I_0), \tag{3}$$

where N is the total number of particles in the pulse, t is the time of arrival, A is the cross-sectional area of the beam at the point of observation, D is the distance between the points of origin and observation for the particles, and L is the total path length of the absorption cell (twice the beam diameter in these experiments). If a relationship between t and δ_{DA} is assumed, t may be eliminated from (3). Under the assumption that the spreading angle α is the same for particles of all velocities, D/t, (3) may be rewritten as

$$\frac{dN}{dv} = -\left(\frac{\pi^{1/2}LD\tan\alpha}{16r_ecf\lambda_0}\right)\ln\left(\frac{I}{I_0}\right), \tag{4}$$

where λ_0 is the wavelength of the emission line. Equation (4) was used to obtain dN/dv from data such as that shown in Fig. 2. An example of the results is shown in Fig. 3. For velocities less than four times thermal (room temperature), (4) has been found to be invalid since the assumption that the emission linewidth is much smaller than δ_{DA} no longer holds and, what is worse, the emission line center appears to be self-reversed, as evidenced by considerably lower sensitivities obtained for transmissivity measurements in cold gas of known densities. As a consequence, dN/dv for the neutrals shown in Fig. 3 is probably a double-peaked distribution with a second peak higher than that shown centered over v_{TH} in spite of the fact that I/I_0 signals received at the corresponding arrival time were buried in the background noise. The area under the portion of the distribution peaked at 2.6×10^3 m/sec represents about 0.25% of the particles injected into the gun.

The Langmuir probe data are also shown in Fig. 2. Since the negatively biased Langmuir probes were centrally located in the intersection of the optical path and the plasma beam, it can be seen that the ions arrive soon after the gun discharge (see third trace from the top) takes place and the neutrals arrive somewhat later. The distance between the gun muzzle and the point of observation was 58 cm for these measurements. From time-of-flight considerations, probe data such as those of Fig. 2 can also be analyzed in terms of a velocity distribution, as shown in Fig. 3. In view of the calibration uncertainties for both the transmittance and Langmuir probe measurements in these initial experiments, the absolute values of the distribution function shown in Fig. 3 are not considered to be more reliable than within an order of magnitude. A pointed indication of this is the fact that if dN/dv for the ions is integrated over v to obtain N, a number is

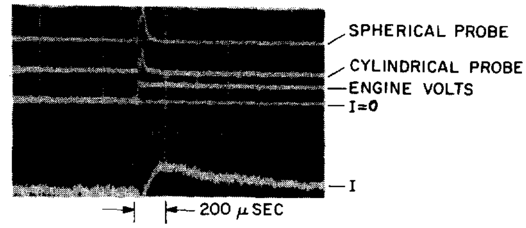

FIG. 2. Langmuir probe signals, gun capacitor voltage, and Xe I 1470-Å transmittance.

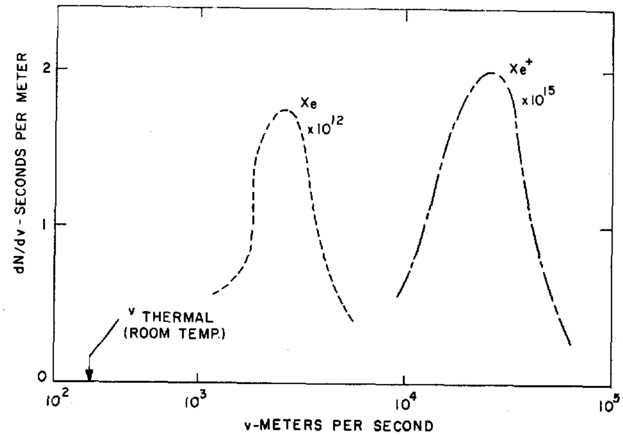

FIG. 3. Velocity distribution function for the particles in the gun exhaust. The peak labeled Xe^+ probably contains other ionic species as well.

obtained which is 2.5 times larger than the number of Xe particles injected into the coaxial gun. This could be accounted for in part by multiply charged xenon ions and impurity ions, all of which have been observed to emit spectral lines.

It has been found that the number density of neutrals is proportional to the mass increment injected into the gun prior to discharge over a range of 20–100 μg per shot. It is especially interesting to note that when less energy per shot was used the neutrals emerged with substantially the same energy as before, but the ions were slower, corresponding to the lower energy. While resonant charge exchange is not ruled out as accounting for the accelerated neutrals observed, the data imply that a peak in the cross section might exist between 1 and 10 eV, which is not borne out in the existing theory.[4]

A spectrogram obtained from a large number of shots revealed the following species to be emitting in the wavelength region between 900 and 2000 Å: Xe I, II, III; C I, II, III; and H. Examining the time profiles of select lines of each of these species revealed that they all occurred in the fast pulse observed with the Langmuir probe, as shown in Fig. 4 for the Xe II line at 1074 Å. This is consistent with some earlier observations made with a multigridded particle probe.[3]

Correlation of these data obtained in the post-

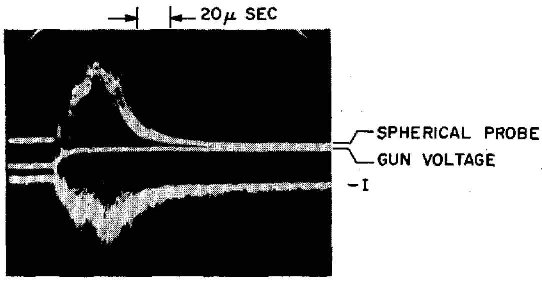

FIG. 4. Spherical Langmuir probe signals (top trace) and XE II 1074-Å emission signal (bottom trace) from the gun exhaust for mass increments of 53 μg and a discharge delay time of 600 μsec. Ten shot overlays are shown.

discharge regime with probe data obtained during the discharge[5] indicates that the neutral component observed may be the result of misalignment of the volume distribution of the current with the predischarge mass distribution.

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Physical Mechanism for the Collisionless Drift Wave Instability

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The growth rate for the collisionless drift instability is calculated in a manner which elucidates the physical mechanism for the instability.

Simple physical models have been useful in understanding fluid instabilities such as the flute instability. However, the collisionless drift wave can be derived only from the Vlasov equation, and its growth rate is determined by resonant particles interacting with the wave. In this note a simple calculation of the growth rate for a drift wave in a nonuniform collisionless plasma immersed in a uniform magnetic field is derived which agrees with the result derived from the Vlasov equation.[1] This simple model has the advantage that the driving mechanism is clear, and it also suggests what effects will be important in determining the growth rate of waves in nonuniform magnetic fields. In general, a wave in the plasma grows when particles transfer energy to the wave, that is,

$$\frac{2\gamma(\boldsymbol{\varepsilon}\cdot\mathbf{E})\cdot\mathbf{E}}{8\pi} = -\mathbf{j}\cdot\mathbf{E}, \tag{1}$$

where E is the perturbed electric field, j is the perturbed particle current, γ is the growth rate, and ϵ is the dielectric tensor.

Consider the plasma configuration shown in Fig. 1 with a drift wave propagating in the electron diamagnetic drift direction. Electrons streaming with their thermal velocity along the field lines interact with the E_z of the drift wave. There are two types of interaction, electrons whose velocity is slightly different than the parallel phase velocity ω/k_z of the wave cause the wave to be Landau-damped but resonant electrons whose velocity is the same as ω/k_z can give energy to the wave. Since the physical mechanism of Landau damping is well known,[2] only the energy transfer by the resonant electrons will be described here. A resonant electron at position a of Fig. 1 experiences a constant $\mathbf{E}\times\mathbf{B}$ drift upward in the x direction. Similarly, at position b of Fig. 1 a resonant electron $\mathbf{E}\times\mathbf{B}$ drifts downward. However, because of the density gradient there are more resonant electrons at position a than at position b. Since the electrons at position a are pushing against the electric field of the wave, there is a net transfer of parallel energy from the resonant electrons to the drift wave. This is the basic driving mechanism for the collisionless drift wave. Now, the magnitude of this effect will be calculated.

The net number of resonant electrons which are moving in the direction of the wave is given by $\delta n_R = \nabla n_R \int (E_y/B)\,dt = \nabla n_R (E_y/B)\tau$ where τ is the time the electrons have experienced a constant E_y. An electron remains resonant with the wave for a time τ if its velocity differs from ω/k_z by less than $\delta V = \lambda_z/2\tau = \pi/(k_z\tau)$. Assuming a Maxwellian distribution in $V_\parallel$, the number of resonant electrons is given by

$$\begin{aligned} n_R &= \frac{n}{(\pi)^{1/2}V_e}\int_{(\omega/k_z)-1/2\delta V}^{(\omega/k_z)+1/2\delta V} \exp\left(-\frac{V^2}{V_e^2}\right) dV \\ &= \frac{n\,\delta V}{(\pi)^{1/2}V_e}\exp\left(-\frac{\omega^2}{k_z^2V_e^2}\right) \\ &= \frac{(\pi)^{1/2}n}{k_zV_e\tau}\exp\left(-\frac{\omega^2}{k_z^2V_e^2}\right). \end{aligned} \tag{2}$$

The energy transferred to the wave by the resonant electrons is given by