Dissociative recombination and the decay of a molecular ultracold plasma

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Abstract.
Double-resonant photoexcitation of nitric oxide in a molecular beam creates a dense ensemble of 51f(2) Rydberg states, which evolves to form a plasma of free electrons trapped in the potential well of an NO+ spacecharge. The plasma travels at the velocity of the molecular beam, and, on passing through a grounded grid, yields an electron time-of-flight signal that gauges the plasma size and quantity of trapped electrons. This plasma expands at a rate that fits with an electron temperature as low as 5 K. Dissociative recombination of NO+ ions with electrons provides the primary dissipation mechanism for the plasma. We have identified three dissociation pathways, and quantified their relative contributions to the measured rate: Two-body dissociative recombination competes with direct three-body recombination to neutral dissociation products, and with a process in which three-body recombination and electron-impact ionization form an equilibrium population of high-Rydberg states that decays by predissociation. Using available collision-theory rate constants for three-body recombination and ionization, together with quantum mechanical estimates of predissociation rates, we predict that the relaxation of the plasma to a high-Rydberg equilibrium outpaces direct three-body dissociative recombination, and, among second-order processes, the rate of two-body electron-cation dissociative recombination substantially exceeds the rate at which the high-Rydberg equilibrium dissociatively relaxes. The rate constant for dissociative recombination extracted from these data conforms with predictions drawn from theory for isolated electron-ion collisions. Methods based on the dissipation of molecular ultracold plasmas may provide a means for estimating rates of dissociative recombination for a variety of complex molecules.

1. Introduction
Ultracold neutral plasmas offer laboratory access to the important limit of a charged particle gas in which the average potential energy of electrostatic interaction approaches the thermal collision energy. Under such conditions, many-body forces govern dynamical properties, and spatial correlations can act significantly to determine the state of a system and its evolution in time.

Most experiments prepare ultracold plasmas from atoms in a magneto-optical trap (MOT). Under typical conditions, radiation from a nanosecond pulsed dye laser excites a substantial fraction of the atoms contained in an initial volume of 1 mm³ at a density of 10¹⁰ cm⁻³.

For an excitation energy tuned just above the ionization threshold, the plasma forms directly. The photon energy in excess of the ionization limit sets the initial electron temperature in this plasma, while the mili-Kelvin trap temperature of the neutral precursor determines that of the ions. Below-threshold excitation to high-Rydberg states forms a Rydberg gas that evolves to
a plasma \( \text{(1) (5) (6) (7)} \). Here, the dynamics of Rydberg-Rydberg and electron-Rydberg collisional ionization processes determine the initial electron temperature \( \text{(8)} \).

Following plasma formation in either case, super-elastic collisions of electrons with Rydberg atoms, initially present or formed by three-body recombination, rapidly raise the electron temperature to 30 K or more. The spatial relaxation of ions moving to positions of lower electrostatic potential energy elevates the ion temperature to the order of 1 K \( \text{(9)} \).

We have recently introduced a new method for the preparation and study of ultracold plasmas \( \text{(10) (11)} \). Our approach begins with a target atom or molecule entrained in a seeded supersonic expansion. Thus far, our experiments have focused on nitric oxide. For NO, seeded at 10 percent in He, expansion cools entrained molecules to a moving frame temperature of about 0.7 K, estimated from a measurement of rotational state populations.

Double-resonant excitation produces a Rydberg gas of \( \approx 5 \times 10^{12} \) molecules \( \text{cm}^{-3} \) in a single selected state. This population evolves to produce prompt free electrons and a durable, cold, quasineutral plasma of electrons and intact \( \text{NO}^+ \) ions. Thereafter, a field of amplitude as low as 3 V/cm – far smaller than that required to field-ionize the precursor state – extracts a small signal of electrons, but pulses of amplitude as high as 200 V/cm fail to destroy the plasma. These observations signal two properties that confirm the evolution to a cold plasma. A low-voltage response establishes the presence of a surface charge of electrons that are trapped by the plasma space charge, but extractable by a weak field. The durability to high fields indicates a Debye screening length much smaller than the diameter of the plasma, shielding the charged particles at its core.

Downstream, transmission through a moveable grid samples the core electrons, profiling their spatial distribution as a function of time. We find that the volume increases at a rate that accelerates with time in accord with the behaviour expected for an electron-driven ambipolar expansion. Analysis of this expansion in the simple limit of the Vlasov equations returns electron temperatures as low as 5 K. Thus, the plasma created in a molecular beam seems both colder and denser than that formed in a MOT, putting this system further into the interesting regime in which the potential and kinetic energies of intermolecular interactions balance, and electrostatic correlations gain in importance.

Significantly, molecular cations carry the positive charge in this \( \text{NO}^+ - e^- \) plasma, which means it can decay by two-body, dissociative recombination. 
Dissociative recombination acts generally to regulate energy balance and composition in conventional plasmas \( \text{(12)} \). Neutralization reactions of \( \text{NO}^+ \) in particular have an important effect on \( \text{NO}_x \) chemistry in upper-atmosphere and combustion contexts \( \text{(13)} \), and have been the subject of significant theoretical interest \( \text{(14) (15)} \). Experimental studies have used both storage-ring and flowing-afterglow techniques to determine accurate thermal rate coefficients. It is interesting now to ask: To what extent can this information predict for the role of dissociative recombination in determining the properties of a nitric oxide molecular ultracold plasma?

Here, we describe techniques by which to monitor dissociative recombination in molecular ultracold plasmas, and characterize conditions under which DR represents the principal channel for decay.

### 2. Experimental Method

Experiments here, in which we study the rate of plasma dissipation, profile the electron density of the plasma volume as it passes through the \((x, y)\) vertical plane of an imaging grid that caps a field-free region. Translation of this grid along the axis \((z)\) of the molecular beam determines the distance traveled by the plasma from its point of preparation to the detection of its electron density. The electron yield detected as a function of flight time gauges the rate at which the plasma decays by recombination of its trapped electrons with \( \text{NO}^+ \) ions.

Figure \[\text{II}\] diagrams the apparatus. A molecular beam, formed by skimming the output of
Figure 1. Experimental apparatus. A skimmed molecular beam passes through grid, G\textsubscript{1}, where it intersects counter-propagating laser beams indicated as \(\omega_1\) and \(\omega_2\). As the excitation volume transits the plane defined by G\textsubscript{2}, a microchannel plate detector (MCP) situated behind G\textsubscript{3} collects the signal of extracted plasma electrons. The components surrounded by the light-grey border translate together on a moving carriage.

The A \(2\Sigma^+\) \(v = 0, J = 1/2\) state decays with a lifetime of 210 ns \cite{16}. Thus, delaying the \(\omega_2\) pulse with respect to \(\omega_1\) yields a reduced number of A-state NO molecules for further excitation. We have no way at present to directly measure the fraction of NO\textsuperscript{*} Rydberg states converted to free NO\textsuperscript{+} ions and electrons, but on thermochemical grounds the observed electron temperature suggests that approximately of 50 percent of NO\textsuperscript{*} - electron collisions drive population upward toward ionization. Simple rate-equation models, employing theoretical rate coefficients for collisional excitation and de-excitation of Rydberg states, support this estimate. We therefore assume that, with zero delay, saturated steps of double resonance ultimately yield a plasma with cations and electrons in approximately equal densities of \(2 \times 10^{12}\) cm\(^{-3}\). Systematically delaying \(\omega_2\) with respect to \(\omega_1\) yields a set of well-defined reduced charged-particle densities.

Our excitation geometry produces a prolate ellipsoid excitation volume with an estimated aspect ratio of 4:1. The long axis of this ellipsoid, which is determined by the width of the molecular beam, lies transverse to the axis of propagation. The short axis conforms with the Gaussian intensity profile of a spatially filtered \(\omega_1\) laser beam, collimated to a full-width at half-maximum of 500 \(\mu m\). The plasma produced in this volume travels with the velocity of the molecular beam over an adjustable distance to pass through the imaging grid, G\textsubscript{2}.

As shown in Figure 1, the carriage supporting this grid, together with G\textsubscript{3} and the MCP
detector, travels on linear bearings supported by four 0.5 inch diameter stainless steel rods. A bellows-isolated motorized actuator controls the position of this carriage to within 0.01 mm over a range of 10 cm.

We generally set $G_2$ to have the same potential as $G_1$ (nominally apparatus ground), creating a field-free region in which we vary the plasma time-of-flight. After $G_2$, the plasma encounters an electrostatic field determined by the potential applied to a third grid, $G_3$, spaced downstream by a fixed distance of 16 mm. For the present experiments, we apply 145 V to $G_3$. As the plasma traverses $G_2$, the field between $G_2$ and $G_3$ extracts electrons which, on the timescale of the measurement, appear instantaneously as signal at the multichannel plate detector. This signal represents the extracted electron density integrated in the transverse dimensions of the plasma, providing a time-dependent trace of this changing density as the plasma volume passes through $G_2$. We assume that the electrons extracted by $G_2$ supply a representative gauge of plasma width and relative charge density.

3. Results

Earlier studies in our laboratory have established that a gas of high-Rydberg NO molecules, entrained in a molecular beam with a longitudinal temperature of ca. 700 mK, evolves to form a quasineutral ultracold plasma of NO$^+$ ions and electrons ($10$, $11$). Present measurements gauge the volume and lifetime of such a plasma, tracking its electron density as a function of time-of-flight over an adjustable interval of distance.

![Figure 2](image)

**Figure 2.** Waveforms captured by the MCP detector for different positions of the moving carriage. Traces have been offset vertically for clarity.

Figure 2 gives a sequence of electron signal traces recorded for an initial Rydberg state in the $nf$ series converging to $N^+ = 2$ for which we have selected the principal quantum number $n = 50$ ($50f(2)$). Successive waveforms show electron signal amplitude as a function of time for increasing $G_2$ displacements. We see that the position of the late signal advances with the position of $G_2$, and that the arrival time of this feature corresponds in each case to the flight time of a neutral volume in the molecular beam from the excitation region to the perpendicular plane of $G_2$. 

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This signal of electrons extracted as the illuminated volume passes through $G_2$ broadens and decreases in amplitude with increasing flight time, which reflects the processes of expansion and decay of the plasma with time. Each trace fits well to a Gaussian function, and we use the parameters of such fits to extract the plasma arrival time, width and relative measure of the total number of electrons. Figure 3 plots the fitted width along the axis of propagation (the short axis of the ellipsoid excitation volume) measured as a function of flight time.

**Figure 3.** Fitted plasma width as a function of centre arrival time at $G_2$. Solid line is a fit of equation 3 to the experimental data, yielding a fitted temperature $T_e + T_I = 5 \text{ K}$.

For any given time, we estimate the total number of electrons remaining in the plasma by integrating the area of the corresponding Gaussian fit. To express this in absolute terms, we use an extrapolation procedure that correlates the earliest-measured data points to the known initial plasma density to calibrate the detector response. For present experiments on the evolution of a Rydberg gas selected to populate the $50f(2)$ state, the signal extrapolated to $t = 0$ for the densest plasma we can create defines the starting point for decay from an estimated initial density $n_e(0) \approx 2 \times 10^{12} \text{ cm}^{-3}$. The introduction of a delay between the $\omega_1$ and $\omega_2$ excitation pulses, initiates decays for plasmas of systematically reduced charged-particle densities. Figure 4 shows a collection of plasma decay waveforms for initial densities varied in this way.

4. Discussion

An ultracold plasma composed of NO$^+$ ions and free electrons expands at a rate dictated by the electron temperature, and decays by dissociative electron-ion recombination. Electron signal, collected as the plasma volume transits the grid $G_2$ provides a measure of its width and relative electron density. By positioning this grid to sample the plasma for varying times of flight, we obtain data that gauge the rates both of plasma expansion and decay.

4.1. Expansion and Electron Temperature

The plasma volume increases as electrons drive radial hydrodynamic ion expansion with a local mean velocity, $\gamma r$. For an ideal, spherical, Gaussian plasma, this velocity grows as:
Figure 4. Integrated area of the plasma signal as a function of arrival time at imaging grid G2 for a range of initial densities. The lines are integrated rate equations fitted to the initial decay data. The calculations consider first-order electron evaporation, plasma expansion and second-order decay processes (c.f. Eqs. 8 and 9)

\[
\frac{d\gamma}{dt} = \frac{k T_e}{m_i \sigma^2},
\]  
where \(\sigma\) is the radius of the plasma and \(T_e\) is the electron temperature. The work of expansion, which converts electron thermal energy into cation hydrodynamic velocity, causes the electron temperature to fall. The initial plasma radius \(\sigma(0)\), together with the electron \((T_e)\) and ion \((T_i)\) temperatures combine to determine a characteristic timescale for both expansion and cooling:

\[
\tau^2 = \frac{m_i \sigma(0)^2}{k_B [T_e(0) + T_i(0)]},
\]
which then yield time-dependent radius and temperature as:

\[
\sigma(t) = \sigma(0) \left[ 1 + t^2 / \tau^2 \right]^{\frac{1}{2}},
\]
and:

\[
T_e(t) = \frac{T_e(0)}{\left[ 1 + t^2 / \tau^2 \right]}.
\]

The present illumination geometry gives our plasma the form of a prolate ellipsoid, as opposed to the self-similar spherical shape described by Eqs.(2) - (4). However, Bergeson and coworkers have obtained fluorescence images of ellipsoidal plasmas showing differential rates of short- and long-axis expansion (17) described by expressions of the form of Eq. (3) for decoupled radial coordinates. Our experiment measures the rate of expansion along the short axis of the ellipsoid.

Assuming that a simple model for ambipolar expansion applies to describe the evolution of the electron density in this coordinate, we fit Gaussian radii, measured as a function of flight time, to Eq. (3), and, in this way, relate a measured rate of plasma expansion to the electron
temperature. We note that the time-dependent width displayed in Figure 3 shows evidence of the acceleration characteristic of an electron-charge-driven ambipolar expansion. A Vlasov fit to this \( \sigma(t) \) returns an initial electron temperature of 5 K.

At the electron density of our plasma, the Wigner-Seitz radius, \( a \), is 500 nm, and such a \( T_e \) would indicate significant correlation (\( \Gamma_e = 7 \), where \( \Gamma_e = q^2/4\pi a \varepsilon_0 k_B T_e \)). The Vlasov equations do not extend to describe the expansion of a charged-particle systems under such conditions. At high correlation, the positive field created by the cations contained in the volume of the plasma lowers the potential energy of free electrons. The depth of this well decreases with expansion, and this elevation of potential energy with expansion acts as a force that suppresses electron pressure. Thus, for the purposes of the following discussion, we can assume that the temperature estimated from rate of plasma expansion represents a lower limit.

### 4.2. Plasma Decay and Dissociative Recombination

The escape of electrons in the first few hundred nanoseconds of plasma formation creates a cationic space charge that acts as an attractive potential for the electrons that remain. Electrons trapped by this potential can decay only by electron-ion recombination. Subsequent plasma expansion, however, acts to reduce the depth of the electrostatic well, and this allows additional electrons to evaporatively escape.

A potential sufficient to trap the plasma electrons under our conditions of \( T_e \) requires an excess cation charge of less than one percent. The decaying plasma thus remains quasi-neutral, and we can represent the mass-action kinetics of charge recombination by differential equations written simply in terms of the time-dependent electron density, \( n_e(t) \).

Elementary processes leading to the irreversible loss of charged particles include electron-electron-ion three-body recombination, followed by predissociation of the formed Rydberg state (collisional dissociative recombination):

\[
\text{NO}^+ + e^- + e^- \xrightarrow{k_{\text{TBR}}} \text{NO}^* + e^- \\
\text{NO}^* \xrightarrow{k_{\text{PD}}} \text{N} + \text{O} \tag{5}
\]

and two-body dissociative recombination:

\[
\text{NO}^+ + e^- \xrightarrow{k_{\text{DR}}} \text{N} + \text{O} \tag{6}
\]

Three-body recombination forms \( \text{NO}^* \) molecules in high-Rydberg states below the ionization threshold of nitric oxide. The rate at which these molecules predissociate determines the kinetics of electron loss by three-body recombination. In the limit where \( k_{\text{PD}} \gg k_{\text{i}o} n_e \), irreversible decay via (3) occurs as a third-order process at the rate of three-body recombination:

\[
- \frac{dn_e(t)}{dt} = k_{\text{TBR}} n_e^2(t) + n_e(t) \frac{dV(t)}{V(t)} dt + k_{\text{ev}} n_e \tag{7}
\]

where we include additional terms to account for the decrease in \( n_e \) owing to plasma expansion and electron evaporation.

In the opposite limit, for \( k_{\text{PD}} \ll k_{\text{i}o} n_e \), decay proceeds via a steady-state population of \( \text{NO}^* \) with a second-order rate law:

\[
- \frac{dn_e(t)}{dt} = \frac{k_{\text{TBR}}}{k_{\text{i}o}} k_{\text{PD}} n_e^2(t) + n_e(t) \frac{dV(t)}{V(t)} dt + k_{\text{ev}} n_e. \tag{8}
\]
Direct dissociative recombination occurs at a rate in accord with its elementary mechanism:

\[- \frac{dn_e(t)}{dt} = k_{DR}n_e^2(t) + \frac{n_e(t)DV(t)}{V(t)} \frac{dV(t)}{dt} + k_{ev}n_e. \tag{9}\]

Considering these elementary mechanisms for irreversible electron loss by two- and three-body recombination, we can expect contributions to the time dependence for plasma decay that, allowing for expansion and electron evaporation, conform with either second- or third-order rate laws. These reactions proceed in parallel to determine the overall rate at which the electron density changes in time. Referring to late electron signal decays observed experimentally, we attempt now to judge the relative importance of each channel.

To predict \(n_e(t)\) according to these rate laws, we represent the plasma volume by a cylinder of concentric shells with radial coordinate, \(r\), reflecting the prolate ellipsoid of ions formed by the intersection of our excitation lasers with the molecular beam. We assign each shell an initial number of electrons according to a radial Gaussian density distribution with a peak density of \(2 \times 10^{12} \text{ cm}^{-3}\), and a radius equal to the value of \(\sigma(0)\) we fit to the data, as shown in Figure 3. Numerical integration of expressions conforming to Eq. 7 or 9 for selected rate coefficients yields the number density within each shell as a function of time.

We vary the volume of each shell according to,

\[ \frac{\partial V}{\partial t} = \frac{\partial V}{\partial r} \frac{\partial r}{\partial t}, \tag{10}\]

using the radial coordinate hydrodynamic mean expansion velocity to define the time dependence of \(r\), viz: \(\partial r/\partial t = \gamma r\), where we determine \(\gamma\) as a function of time from the fitted value of \(\tau\). Integrating the calculated number of electrons over the volume of the model gives the total number of plasma electrons as a function of time, which we compare with the integrated area of the plasma signal.

The rate of electron loss proceeds at third-order in charged particles only for the mechanism of three-body recombination in the limit for which \(k_{PD} \gg k_{ion}n_e\). In a recombining atomic ultracold plasma, the balance between three-body recombination and electron-impact ionization creates a Saha equilibrium (18) between free charges and high-Rydberg states that populates a band of principal quantum numbers between a kinetic bottle neck at \(n = \sqrt{R_g/4k_BT_e}\) and the thermochemical limit for bound states at \(n = \sqrt{R_g/k_BT_e}\) (19). Reliable models enable the calculation of temperature-dependent electron-impact ionization rate coefficients in this range (8), on the basis of which our conditions yield an initial \(k_{ion}n_e\) product of \(4.6 \times 10^9 \text{ s}^{-1}\).

Remacle and Vrakking (20) have modelled predissociation in optically populated NO Rydberg states with principal quantum numbers near 100. They find that fields as small as 40 mV cm\(^{-1}\) serve to mix fast-predissociating low-\(l\) zero-order states with non-penetrating, higher-\(l\) states of the same and different principal quantum numbers. This dilution of coupling strength to neutral decay channels produces biexponential decays with a short-component time constant on the order of 400 ns and a long-component decay time that exceeds 12 \(\mu\)s for \(n = 105\). In contrast to optical excitation, three-body recombination directly addresses orbitals of high-angular momentum. For this reason, we can expect \(k_{PD} \approx 1 \times 10^5 \text{ s}^{-1}\) for a collisionally populated Rydberg state with \(n \approx 100\), falling as \(n^{-3}\) to \(1 \times 10^{-4} \text{ s}^{-1}\) at \(n \approx 180\), resulting in a relaxation rate to Saha equilibrium over this interval that far outpaces predissociation. Thus, we should not expect to observe a rate of decay that displays a third order dependence on the density of charged particles, and fits to the experimental data readily confirm this.

Solid lines through the experimental decays in Figure 4 represent fits to generalized second-order rate expressions of the form of Eqs. (8) and (9) for a set of systematically varied initial charged-particle densities. These fits yield a second-order electron-ion recombination
rate constant of $k_{\text{eff}} = 7 \pm 5 \times 10^{-6}$ cm$^3$ s$^{-1}$, and first-order electron evaporation rate of $k_{\text{ev}} = 1 \pm 0.4 \times 10^3$ s$^{-1}$. This phenomenological second-order rate constant, $k_{\text{eff}}$, contains contributions from both dissociation mechanisms (5) and (6):  

$$k_{\text{eff}} = \frac{k_{\text{TBR}}}{k_{\text{ion}}} k_{\text{PD}} + k_{\text{DR}}. \quad (11)$$

It remains now to decide which of these terms is more important, whether the rate process associated with electron loss is driven by three-body recombination accompanied by predissociation, or represents a conventional, two-body NO$^+$– electron recombination leading directly to neutral dissociation products.

To judge the influence of three-body recombination, we return to the population of high-Rydberg states in the principal quantum number range between $\sqrt{R_y/4k_B T_e}$ and $\sqrt{R_y/k_B T_e}$, which is sustained by the Saha equilibrium between collisional recombination and ionization. The corresponding equilibrium constant, which equates to the ratio of forward to reverse rate constants, is determined by statistical factors consisting simply of the density of bound Rydberg states and the partition function for free electrons:

$$\frac{k_{\text{TBR}}}{k_{\text{ion}}} = \frac{n_i^2 \Lambda^3}{\exp(-R_y/n_i^2 k_B T_e)} \quad (12)$$

where $\Lambda = \sqrt{h^2/2\pi m_e k_B T_e}$ is the de Broglie wavelength for electrons with temperature $T_e$.

Over the relevant range of high-Rydberg principal quantum numbers, $n_i$, from approximately 90 to 180, this ratio varies slowly from 1 to $10^{-11}$ to $3 \times 10^{-12}$ cm$^3$. Assuming a statistical distribution of population over these high Rydberg states and scaling the $n = 105$ predissociation rate by $n_i^{-3}$, we obtain an average value of the rate constant for collisional dissociative recombination, $(k_{\text{TBR}}/k_{\text{ion}})k_{\text{PD}}$, of $6 \times 10^{-8}$ cm$^3$ s$^{-1}$. Thus, we can conclude that three-body recombination makes a minimal contribution to irreversible electron loss in this molecular ultracold plasma.

Experiments in afterglows and accelerators have measured rate constants for dissociative recombination of NO$^+$ at temperatures as low as 200 K [21] [22] [23]. These observations conform well with ab initio R-matrix – multichannel quantum defect theory calculations of $k_{\text{DR}}$ for electron temperatures from 1 to 5000 K [21]. In the low-temperature range from 5 to 10 K, theory predicts a two-body rate constant for dissociative recombination from 2 to $3 \times 10^{-6}$ cm$^3$ s$^{-1}$. Thus, even though the present conditions of electron correlation limit the certainty with which the observed expansion rate determines the plasma electron temperature, the corresponding variation in theoretically predicted rate constant does not exceed the present statistical uncertainty in our measurement. Perhaps with improved accuracy, determinations of the dissociative recombination rate can serve as a probe of plasma conditions.

Based on results obtained thus far, we can draw the following conclusions concerning the dissipation of a molecular ultracold plasma of nitric oxide. The decay of plasma density does not proceed as a rate process that is third-order in charged particles. Also, the population of high-Rydberg states maintained in equilibrium with ions by three-body recombination does not play a large role in determining the second-order decay rate. Instead, it appears that irreversible charged-particle loss occurs in simple two-body collisions by conventional electron-ion dissociative recombination. Even though this plasma seems demonstrably to be at or near conditions for Coulomb correlation, rate processes do not appear strongly perturbed by many-body non-ideal collisional effects. The rate constant for dissociative recombination extracted from these data conforms with predictions drawn from theory for isolated electron-ion collisions. Methods based on the dissipation of molecular ultracold plasmas therefore may eventually offer a practical means for measuring rates of dissociative recombination for a variety of complex molecules.
5. Acknowledgements
It is a pleasure to acknowledge helpful discussions with T. Pohl and J. M. Rost. This work was supported by the Natural Sciences and Engineering Research Council of Canada (NSERC), the Canada Foundation for Innovation (CFI) and the British Columbia Knowledge Development Fund (BCKDF).

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CORRIGENDUM

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2011 Journal of Physics: Conference Series 300 012005

Received: 25 July 2011
Published: 27 July 2011

Figures 3 and 4 were corrupted in the PDF conversion of the original manuscript. This changed the coordinates of the tics marking the upper and lower limits of the error bars, producing what look like additional data. The plots below show the correct versions of Figures 3 and 4.

Figure 3. Fitted plasma width as a function of centre arrival time at G2. Solid line is a fit of equation 3 to the experimental data, yielding a fitted temperature $T_e + T_I = 5$ K.

![Figure 3](image-url)
Figure 4. Integrated area of the plasma signal as a function of arrival time at imaging grid G2 for a range of initial densities. The lines are integrated rate equations fitted to the initial decay data. The calculations consider first-order electron evaporation, plasma expansion and second-order decay processes (c.f. Eqs. 8 and 9).