STORAGE RING MEASUREMENT OF THE C IV RECOMBINATION RATE COEFFICIENT

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ABSTRACT

The low-energy C IV dielectronic recombination (DR) rate coefficient associated with $2s \rightarrow 2p$ $\Delta n = 0$ excitations of this lithium-like ion has been measured with high-energy resolution at the heavy-ion storage ring TSR of the Max-Planck-Institut für Kernphysik in Heidelberg, Germany. The experimental procedure and especially the experimental detection probabilities for the high Rydberg states produced by the recombination of this ion are discussed in detail. From the experimental data a Maxwellian plasma rate coefficient is derived with $\pm 15\%$ systematic uncertainty and parameterized for ready use in plasma-modeling codes. Our experimental result especially benchmarks the plasma rate coefficient below $10^4$ K where DR occurs predominantly via C III ($1s^22p^4l$) intermediate states and where existing theories differ by orders of magnitude. Furthermore, we find that, to within our systematic uncertainty of 15%, the total dielectronic and radiative C IV recombination can be represented by the incoherent sum of our DR rate coefficient and the radiative recombination rate coefficient of Pequignot and coworkers.

Subject headings: atomic data — atomic processes — X-rays: general

1. INTRODUCTION

Carbon is one of the cosmically most abundant elements. Consequently, line emission from carbon ions is observed from a wide range of cosmic objects. In low-density, photoionized and electron-ionized cosmic plasmas, the dominant mechanisms for recombination are radiative recombination (RR) and dielectronic recombination (DR). For C IV these processes can be expressed as

$$C^3\,^+(1s^22s) + e^- \rightarrow C^2\,^+(1s^22nl) + h\nu \quad (1)$$

and

$$C^3\,^+(1s^22s) + e^- \rightarrow C^2\,^+(1s^22npl) \rightarrow \begin{cases} C^2\,^+(1s^22nl) + h\nu, \\ C^2\,^+(1s^22pn'l') + h\nu, \end{cases} \quad (2)$$

respectively. In low-energy C IV DR, the $1s^22s$ ion is excited by a $2s \rightarrow 2p$ $\Delta n = 0$ transition (excitation energy $\sim 8$ eV) only within the $n = 2$ shell. At higher energies above 37.55 eV, the $2s$ electron can be excited to $3l$/substates ($\Delta n = 1$). The lowest energy resonances of the associated DR series with configurations $1s^23l'$ of C III are at energies above about 12 eV. The strengths of these resonances can be expected to be very much smaller than those associated with $2s \rightarrow 2p$ core transitions. The K shell cannot be excited with energies smaller than 240 eV. In this paper we exclusively deal with $\Delta n = 0$ DR.

The calculation of DR rate coefficients is a challenging task since an infinite number of states are involved in this process. Moreover, relativistic and many-body effects have to be accounted for in high orders even in the case of DR of a light ion such as C IV (Mannervik et al. 1998). Approximations and computational simplifications are needed in order to make DR calculations tractable. It turns out that different calculations yield rate coefficients differing by up to orders of magnitude. In this situation benchmarking experiments are vitally needed in order to guide the development of the theoretical methods and to provide reliable DR rate coefficients for plasma modelers. For C IV the available theoretical rate coefficients have recently been critically compared by Savin (2000).

In the past decade electron coolers at heavy-ion storage rings have developed into the most successful experimental tool for electron-ion recombination studies (Schuch 1993; Müller & Wolf 1997). Currently, corresponding research programs are carried out at the heavy-ion storage rings ESR of the Gesellschaft für Schwerionenforschung (GSI) in Darmstadt (Brandau et al. 1997), TSR of the Max-Planck-Institut für Kernphysik in Heidelberg (Müller et al. 1997, 1998; Wolf et al. 2000) and CRYRING of the Manne-Siegellabor-Laboratory in Stockholm (Schuch et al. 1997, 1998). Extensive bibliographic compilations on DR measurements at storage rings have been published, e.g., by Müller (1995) and Schippers (1999). The basic approach for deriving plasma rate coefficients from storage ring measurements has been summarized by Müller (1999). Recent experimental work on plasma rate coefficients for astrophysical and other plasma applications has been published by Savin et al. (1997, 1999) on DR of Fe xviii and Fe xix, by Schippers et al. (1998) on DR of Ti v, and by Schippers et al. (2000) on DR of Ni xxvi.

In this paper we present the measured C IV recombination rate coefficient. The paper is organized as follows: In § 2 the experimental procedure is described. Particular attention is given to the field ionization of high-$n$ Rydberg states in motional electric fields that are unavoidable with the experimental arrangement at an ion storage ring. In § 3 the experimental result is presented and the impact of field ionization on the measured DR resonance strength is discussed. A theoretical estimate of the unmeasured DR rate is presented, and the C IV plasma DR rate coefficient is derived. In § 4 our derived rate is compared to existing theoretical results. Finally, in § 5 we derive the total DR + RR C IV rate coefficient and compare it to a unified calculation of Nahar & Pradhan (1997). A model calcu-
2. EXPERIMENT

The C IV recombination measurements have been performed at the heavy-ion storage ring TSR (Jaeschke et al. 1989) of the Max-Planck-Institut für Kernphysik in Heidelberg, basically following the procedure of earlier measurements (Kilgus et al. 1992; Lampert et al. 1996). A beam of $^{12}$C$^{+}$ ions at an ion energy $E_i \approx 1.5$ MeV $\cdot$ u$^{-1}$ was supplied by the MPI accelerator facility and injected into the TSR. In the TSR electron cooler, situated in one of the straight sections of the storage ring (see Fig. 1), the circulating C$^{+}$ ions were merged with a beam of electrons moving collinearly with the ions at roughly identical velocity. At matched beam velocities the electron beam had an energy $E_e = (m_e/m_i)E_i \approx 840$ eV. It was merged with the ion beam over a length of 1.5 m and was guided by a longitudinal magnetic field of 42 mT over its entire path. The electron density was $(5-7) \times 10^6$ cm$^{-3}$ depending on the laboratory electron energy $E_e$. The electron beam’s diameter was 3 cm. Electron cooling, i.e., the ion interaction with the overlapping electron beam, on a timescale of $\approx 1$ s, compresses the circulating ion beam to a diameter of 1–2 mm; simultaneously, the longitudinal velocity of the freely coating ions adjusts itself to the average electron velocity. By repeated injections, while keeping the cooled beam stored, C$^{+}$ ions were accumulated (Grieser et al. 1991) on a timescale of $\approx 60$ s up to electrical currents of typically 30 $\mu$A, corresponding to $\approx 8 \times 10^7$ stored ions.

After ion accumulation, recombination rates of C$^{+}$ ions were measured by counting recombed $^{2}$C$^{+}$ ions on a multichannel-plate detector located behind the first dipole magnet (bending radius 115 cm) downstream of the electron cooler (see Fig. 1). The dipole magnet keeps the circulating C$^{+}$ ion beam on a closed orbit while it deflects the recombined $^{2}$C$^{+}$ ions less strongly so that they hit the detector. Absolute rate coefficients for the recombination of C$^{+}$ ions with electrons in the collinear overlap region were measured as a function of the average relative energy

$$\hat{E} \approx (\sqrt{E_e} - \sqrt{E_i})^2$$

between electron and ion beams. The relativistically correct expression (see, e.g., Schippers et al. 2000) has been used in the data analysis. The laboratory electron energy $E_e$ was deduced from the cathode voltage and electron current applying a correction for the electron-space charge (Kilgus et al. 1992).

The electron motion in the transverse degrees of freedom, confined by the magnetic field, is largely decoupled from that in longitudinal direction. In a comoving reference frame, the electron velocity spread is characterized by the longitudinal (with respect to the electron beam direction) and transverse temperatures $T_\parallel$ and $T_\perp$, with $k_B T_\parallel \approx 10$ meV and $k_B T_\perp \approx 0.15$ meV $\ll k_B T_i$ ($k_B$ denotes the Boltzmann constant). Both temperatures are considerably lower than the cathode temperature, the low longitudinal temperature resulting from the electron acceleration and the low transverse temperature resulting from adiabatic magnetic expansion of the electron beam (Pastuszka et al. 1996).

The electron velocity distribution is represented by

$$f(\vec{v}, \vec{v}) = \frac{m_e}{\sqrt{2\pi k_B T_\parallel}} \exp\left[-\frac{m_e(\vec{v}_i - \vec{v})^2}{2k_B T_\parallel}\right] \times \frac{m_e}{2\pi k_B T_\perp} \exp\left[-\frac{m_e v_z^2}{2k_B T_\perp}\right].$$

The experimental energy resolution corresponds to the width of this distribution and amounts to (Müller 1999)

$$\Delta E(\text{FWHM}) = \left\{[\ln(2) k_B T_\parallel]^2 + 16\ln(2) k_B T_\parallel \hat{E}^2\right\}^{1/2}.$$ 

As has been demonstrated by Bartsch et al. (1999, 2000) and Schippers et al. (2000), DR rate coefficients can considerably be influenced by external electric and additional crossed magnetic fields in the TSR electron cooler. Therefore, care has been taken to minimize such fields in the interaction region. In addition to the magnetic field of 42 mT, only small electric stray fields are present in the interaction region. Components of the magnetic field transverse to the ion velocity produce motional electric fields in the ion rest frame amounting to 2 V cm$^{-1}$ for the estimated maximum field angles of about 0.3 mrad. Small stray fields are also expected from the electron space charge (Kilgus et al. 1992). By monitoring the velocity of the stored, electron-cooled ion beam, the ions are centered within the electron beam to about $\pm 2$ mm. For aligned electron and ion beams, the electric space-charge field remains below 1 V cm$^{-1}$ over the ion beam cross section of $\leq 2$ mm diameter. Altogether, electric stray fields below 3 V cm$^{-1}$ are estimated. Considering the findings by Bartsch et al. (1999, 2000) and Schippers et al. (2000), we conclude that a reasonably “field-free” DR measurement can be carried out under these conditions.

Before their detection, the recombined $^{2}$C$^{+}$ ions have to travel through the toroidal magnet that guides the electron beam out of overlap, through correction dipoles and quadrupoles, and through the deflection dipole that separates the electron beam from the ion beam (Fig. 1). All these components cause transverse magnetic fields considerably higher than those in the interaction region, leading to motional electric fields that can field ionize recombined ions in highly excited states. Since a sizable fraction of recom-

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**Fig. 1.** Sketch of the experimental setup. The C$^{+}$ ion beam enters the electron cooler from the right. After traveling through the merging section inside the cooler, the ions pass the toroidal magnet used for deflecting the electron beam out of the ion beam’s path, two correction dipole magnets, and the charge analyzing dipole magnet. The respective flight times, $t^{(1)}_F = 55$ ns, $t^{(2)}_F = 103$ ns, $t^{(3)}_F = 133$ ns, and $t^{(4)}_F = 275$ ns, from the center of the electron cooler to these magnets are indicated; $t_l = 87$ ns is the flight time through the merging section. Recombined C$^{+}$ ions are counted with the recombination detector. Not shown are correction magnets to the right of the cooler and focusing magnets in between the correction and charge-analyzing dipoles.
combined ions are expected to be formed in highly excited states, the efficiency of their detection needs particular consideration. Recombined ions that reach a zone of large motional electric fields in sufficiently excited states will be field ionized and hence cannot be detected; however, ions formed in highly excited states that are able to decay radiatively to lower levels before reaching the critical field region will still be detected. At the given measuring conditions, the 1.5 m nominal length of the interaction region corresponds to a flight time of 87 ns. The motional electric field rises to ≈4 kV cm⁻¹ in the toroid after a flight time of 55 ± 44 ns, depending on the location along the interaction path where the recombination takes place. Two correction magnets with peak fields of ~6 and ~12 kV cm⁻¹ are reached after 103 ± 44 and 133 ± 44 ns, respectively, and the deflection dipole with a peak field of ≈106 kV cm⁻¹ after 275 ± 44 ns. The critical quantum number for field ionization of an ion in an electric field \( F \) is estimated from (Gallagher 1994)

\[
n_F = 4 \sqrt{\frac{q^3}{9F}},
\]

where \( q \) is the charge of the ion core and \( F \) is measured in atomic units (1 au = 5.142 × 10⁻¹⁹ V cm⁻¹). For \( \text{C}^2⁺(nl) \) the core charge state is \( q = 3 \), and thus the critical quantum number is \( n_F = 19 \) for \( F = 106 \) kV cm⁻¹. On the other hand, the radiative lifetime of the 19p level in \( \text{C}^2⁺ \) is 16 ns (hydrogenic value) and hence considerably shorter than the 275 ns flight time from the center of the cooler to the charge analyzing dipole magnet. The detection probability for recombination events leading to high Rydberg levels \( n \geq 20 \) will therefore depend on the competition between the radiative decay and the flight time to regions where motional electric fields are present. Reciprocally, the experimental setup offers a detection probability near unity for all recombination events leading to states with principal quantum numbers at least up to about 20.

With a cooled and accumulated beam of \( \text{C}^3⁺ \) ions of typically 30 μA freely circulating in the ring, measurements of the recombination rate were performed by stepping the acceleration voltage in the electron cooler to values different from those at cooling for short time intervals and recording the detector count rate after the acceleration voltage had settled to a stable value. In sequence, the acceleration voltage was set to the value required to obtain the desired relative energy where the recombination rate coefficient was to be measured and then to a value producing a fixed, high relative energy where the count rate of \( \text{C}^2⁺ \) ions was dominated by capture in the residual gas and where a background rate was measured. The corresponding time intervals of typically 20 ms duration will be called “signal” and “reference” windows, respectively. In between the voltage jump and the start of data taking, a settling time of 1.5 ms allowed the power supplies to reach the preset acceleration voltage. After the signal and reference intervals, the acceleration voltage was stepped back to the cooling value and kept there for 30 ms to ensure a fixed ion energy \( E_i \), and a well-cooled ion beam for each measurement. Measurement cycles were repeated scanning the acceleration voltage level applied during the signal window. A scan comprising 420 data points took ~30 s, and during that time the stored \( \text{C}^3⁺ \) ion current decreased only slightly; the ion current was restored by “stacking” a few new injections, and several such scans were repeated until the desired integral ion counts were collected for the recombination spectrum. Our experimental range of relative energies (0–10.5 eV) was covered by three overlapping scans.

The raw data consist of time-correlated signal and reference count rates \( R'(E_s) \) and \( R'(E_r) \) for each value of the laboratory electron energy \( E_s \) reached during the scans. Both rates include the background from electron capture of \( \text{C}^3⁺ \) ions in the residual gas, and since they are measured in time windows only milliseconds apart, the background is expected to cancel in the difference \( R'(E_s) - R'(E_r) \) even in the case of variations of the residual gas pressure on a timescale as fast as seconds. After conversion to the scale of relative energies, \( \tilde{E} \), the energy-dependent recombination rate coefficient \( \alpha(\tilde{E}) \) is obtained through

\[
\alpha(\tilde{E}) = \frac{(R'(\tilde{E}) - R'(\tilde{E}))^2}{n_s N_i (L/C) \eta} + \alpha(\tilde{E}) \frac{n_r}{n_s}, \tag{7}
\]

where \( n_s \) and \( N_i \) denote the electron density and the number of stored ions, respectively, \( L = 1.5 \) m is the nominal length of the interaction region, and \( C = 55.4 \) m is the circumference of the storage ring. At the relatively high experimental ion energy, the efficiency of the channel-plate recombination detector can be assumed to be \( \eta = 1 \). The factor \( \gamma \) accounts for the relativistic transformation between laboratory and center-of-mass frames. The second term in equation (7) is a small correction that re-adds the electron-ion recombination rate at the reference point. With a suitably chosen reference point—\( \tilde{E} = 10.5 \) eV in the present case—this rate is due to RR only. Here it is calculated using a modified semiclassical formula for the RR cross section (see eq. [13] below) to be \( \alpha(\tilde{E}) \approx 3.5 \times 10^{-13} \) cm³ s⁻¹. Since this is only a small correction to the experimental data at lower energies, the insertion of a rough theory value is justified.

In the toroidal magnetic sections of the electron cooler, the electron beam was guided in and out of the interaction region on two 45° bends with a radius of 80 cm (see Fig. 1). In the first section the ion beam interacts with the electron beam at decreasing and in the second section at increasing angles between the two beams. This leads to shifted relative energies over regions of \( \approx 20 \) cm length each. From the known geometry, corrections for these merging regions are included in the data evaluation (Lampert et al. 1996).

The experimental collision energy (\( \tilde{E} \)) scale is known under the conditions of the present measurement within a systematic relative error of about 1%; even higher energy accuracies were reached after a precise adjustment of the energy scale to known spectroscopic Rydberg series limits (see below). The systematic uncertainty in the absolute recombination rate coefficient is estimated to be \( \pm 15\% \), where the dominant errors are due to the ion and electron current determination and the detection efficiency. The influence of the inaccurate knowledge of the effective overlap length \( L \) is strongly reduced by the toroid correction (Lampert et al. 1996). The statistical uncertainty of the results presented below amounts to less than 1% in the rate coefficient maximum.

3. RESULTS

Our experimental \( \text{C} \ IV \) rate coefficient comprising RR and DR contributions is displayed in Figure 2. The RR peak at \( \tilde{E} = 0 \) eV and individual members of the \( 2\text{nl} \)
Rydberg series of DR resonances are resolved for 4 ≤ n ≤ 12. As the 2p_{1/2}-2p_{3/2} splitting in C IV amounts to only 0.013 eV (Edlén 1983), it is not possible to observe two separate Rydberg series with the given experimental resolution, in contrast to measurements on heavier lithium-like ions (Kilgus et al. 1992; Schippers et al. 2000; Bartsch et al. 2006). The series limit E∞ corresponding to the 2s → 2p excitation energy was determined by extrapolating the observed Rydberg resonance positions

\[ E_n = E_\infty - \mathcal{A} \frac{q^2}{n^2} \]  

from n = 5 – 12 to infinite n. In equation (8) the ion charge q = 3 and \( \mathcal{A} = 13.6056 \) eV have been inserted. In Figure 2 and throughout the paper, we have adjusted the collision energy (\( \bar{E} \)) scale by multiplying it with a constant factor differing from unity by 0.65%\( \gamma_0 \), which brings the extrapolated series limit into agreement with the spectroscopic value of 8.005 eV (Edlén 1983).

In order to obtain the C IV DR rate coefficient, the RR contribution has been removed from the spectrum of Figure 2 by subtracting an empirical function \( \sigma_{\text{RR}}(\bar{E}) = a_0 + a_1 \bar{E} + a_2(1 + a_3 \bar{E} + a_4 \bar{E}^2) \) from the experimental rate coefficient, with the coefficients \( a_i \) determined by fitting \( \sigma_{\text{RR}}(\bar{E}) \) to those parts of the spectrum that do not exhibit DR resonances. The corresponding energy intervals that have been used in the fit were 0.021 – 0.1 eV, 1.0 – 2.0 eV, and 8.5 – 10.5 eV. The parameters obtained from the fit are \( a_0 = 1.556 \times 10^{-12} \) cm³ s⁻¹, \( a_1 = -1.929 \times 10^{-13} \) cm³ s⁻¹ eV⁻¹, \( a_2 = 4.726 \times 10^{-11} \) cm³ s⁻¹, \( a_3 = 20.89 \) eV⁻¹, and \( a_4 = -3.665 \times 10^{-5} \) eV⁻². With these parameters, \( \sigma_{\text{RR}}(\bar{E}) \) is defined in the range 0.0021 eV ≤ \( \bar{E} \) ≤ 10.5 eV.

In the inset of Figure 2, we compare the part of the DR spectrum that comprises the 2p_{1/2} DR resonances with the measurement of Mannervik et al. (1998) carried out at the heavy-ion storage ring CRYRING in Stockholm. The published C IV CRYRING data extend only to \( \bar{E} = 0.7 \) eV. For the comparison we have removed the RR background from the present and from the CRYRING measurement by using the procedure described above. The comparison shows that the two measurements have a similar energy resolution. The overall shape of the spectrum and the peak positions are almost the same. However, when integrating over the 0.1 – 0.7 eV energy range, we obtain for the 2p_{1/2} manifold the resonance strength \( 1.9 \times 10^{-11} \) and \( 2.5 \times 10^{-11} \) cm³ s⁻¹ from the CRYRING and from the present TSR measurements, respectively. This difference is within the 30% summed uncertainty (15% for each experiment individually, assuming a systematic uncertainty for the CRYRING measurement similar to that for the present one) for the absolute value of the measured rate coefficient. Our value is somewhat closer to the theory value \( 3.3 \times 10^{-11} \) cm³ s⁻¹ of Mannervik et al. (1998).

For the derivation of a meaningful plasma rate coefficient from our experimental data, we have to estimate how much DR strength is not measured due to the cutoff of high Rydberg states in the electric fields on the recombinated ions’ path from the cooler to the detector. To this end we have carried out C IV DR calculations using the atomic structure code AUTOSTRUCTURE (Badnell 1986). In order to be able to compare the calculated DR cross section \( \sigma \) with our measured rate coefficient, we have performed the convolution

\[ \sigma(\bar{E}) = \sum_n \sigma_{nl}(\bar{E}v)w(\bar{E}, v)dv, \]

where we have used the electron velocity distribution of equation (4) with \( k_b \) Tg = 0.15 meV and \( k_B \) Tl = 10 meV. In the convolution, energy average cross sections (energy bin width 9.52 meV) have been treated as delta functions. In equation (9), \( \sigma_{nl} \) is the nl-specific detection probability of a recombinated ion, with the outer electron being in a Rydberg state characterized by the quantum numbers n and l. Correspondingly, \( \sigma_{nl} \) denotes the cross section for DR via this Rydberg state. When comparing the resonance strengths of the individually resolved 2pnl DR resonances with 5 ≤ n ≤ 10, we find that the calculation yields values that are on average a factor 1.25 higher than the experimental ones. Therefore, we have scaled down the calculation by a factor 0.8. Additionally, the theoretical energy scale has been adjusted by shifting it by 0.06 eV toward higher energies such that the spectroscopic value for the series limit is reproduced.

The comparison between our scaled AUTOSTRUCTURE calculation and our experimental result is shown in Figure 3. The different curves correspond to different assumptions for the detection probabilities of high Rydberg states. The solid line has been obtained by setting \( \gamma_{nl} = 1 \) for all nl up to \( n = 1000 \), which is the maximum n used in the calculation. By inclusion of such high-n states, we have made sure that the calculation has converged; i.e., calculations up to \( n = 900 \) and \( n = 1000 \) yield rate coefficients that are indistinguishable from one another. It is obvious that the larger fraction of the resonance strength due to DR via high Rydberg states has not been measured. It should be noted that this effect becomes much less significant for DR on more highly charged ions where comparatively less DR strength is accumulated in high-n DR resonances (Savin et al. 1997, 1999; Schippers et al. 1998, 2000). Moreover, Rydberg electrons are more tightly bound in highly charged ions and hence require stronger fields to become ionized in the charge-analyzing field of the apparatus.
The simple picture of a hard cutoff at \( n = n_p = 19 \), as derived in § 2, is not a good description of the field-ionizing properties of our apparatus. This can be seen from the comparison of the experimental data with the dashed curve in Figure 3, which has been obtained with \( Y_{nl} = 1 \) for \( n \leq 19 \) and \( Y_{nl} = 0 \) for \( n > 19 \). A detailed model that takes into account radiative decay of higher Rydberg states on the way from the cooler to the field ionization magnet is described in Appendix A. The resulting calculated \( nl \)-specific detection probabilities \( Y_{nl} \) are plotted in Figure 4. Via equation (9) they yield the dash-dotted curve in Figure 3. Apparently, the contributions from Rydberg states with \( n > 19 \) to the measured rate coefficient are not negligible. Figure 4 indicates that the main contribution from these high-\( n \) Rydberg manifolds is by the short-lived \( p \) states. On the other hand, the cutoff of Rydberg states with \( n > 45 \) in the toroid (see Table 2) has a decisive influence on the detection probabilities since the time available for the radiative decay of higher \( n \) states to below \( n = 45 \) is too short (see Appendix A). The remaining discrepancy between the outcome of the detailed model and the measured DR spectrum for very high Rydberg states is attributed to model-inherent simplifications. For example, very high \( n \) states are easily perturbed by even small stray fields, possibly leading to enhanced radiative decay rates for \( l \neq 1 \) states due to \( l \)-mixing within \( n \) manifolds, especially in the field-ionization regions. Such effects are not taken care of by the model.

Finally, we note that at low energies the AUTOSTRUCTURE calculations do not reproduce the measured 2\( p4l \) DR resonance structure (Fig. 3, inset), a result that had to be expected given the theoretical effort described by Mannervik et al. (1998) to be necessary for matching the experiment (Fig. 2, inset).

In view of the substantial experimental cutoff of high Rydberg states, we derive our C\textsc{iv} \( \Delta n = 0 \) DR rate coefficient in a plasma by using the experimental DR spectrum only below \( \bar{E} = 7.6 \) eV. Above that energy we substitute the experimental DR spectrum by the scaled AUTOSTRUCTURE result with a cutoff at \( n = 1000 \) (Fig. 3, solid line). The composite DR spectrum is then converted into a cross section \( \sigma(\bar{E}) = \sigma(\bar{E})/\sqrt{2\bar{E}/m_e} \) and convoluted with an isotropic Maxwellian electron energy distribution, yielding the plasma rate coefficient

\[
\sigma(T_e) = (k_B T_e)^{-3/2} \frac{4}{\sqrt{2m_e \pi}} \int_0^\infty d\bar{E} \sigma(\bar{E}) \bar{E} \exp \left( \frac{-\bar{E}}{k_B T_e} \right) \] (10a)

\[
= (k_B T_e)^{-3/2} \frac{2}{\sqrt{\pi}} \int_0^\infty d\bar{E} \sigma(\bar{E}) \bar{E}^{1/2} \exp \left( \frac{-\bar{E}}{k_B T_e} \right) \] (10b)

at the plasma electron temperature \( T_e \). This procedure is safe as long as the relative energy \( \bar{E} \) is larger than the experimental energy spread defined in equation (5), i.e., for \( T_e \gg T \approx 120 \) K (see discussion below).

The resulting plasma C\textsc{iv} DR rate coefficient is displayed as the thick solid line in Figure 5. The curve displays two local maxima. The first one, which is due to the 2\( p4l \) DR resonances, is especially benchmarked by our experiment. The second one is caused by DR via high-\( n \) Rydberg states. Here our result is dominated by the AUTOSTRUCTURE calculation that has been adjusted by a constant factor 0.8 to our experiment. For comparison, the dotted curve in Figure 5 represents the plasma rate coefficient that has been derived directly from the measured DR spectrum (Fig. 3, filled symbols) without the additional resonance strength above 7.6 eV (Fig. 3, shaded area) introduced via the AUTOSTRUCTURE calculation. At temperatures above 10,000 K it is about a factor 5 lower than the composite rate coefficient. On the other hand, the calculated part of the
composite DR spectrum does not influence our result at temperatures below 10,000 K.

A convenient representation of the plasma DR rate coefficient is provided by the following fit formula:

\[ \alpha(T_e) = \frac{T_e}{2} \sum_i c_i \exp \left(-\frac{E_i}{k_B T_e}\right). \]  

(11)

It has the same functional dependence on the plasma electron temperature as the Burgess (1965) formula, where the coefficients \( c_i \) and \( E_i \) are related to oscillator strengths and excitation energies, respectively. The results for the fit to the experimental C IV \( \Delta n = 0 \) DR rate coefficient in a plasma are summarized in Table 1. The fit deviates from the thick solid line in Figure 5 by no more than 1% for 650 K < \( T_e < 1500 \) K and by no more than 0.2% for \( T_e \geq 1500 \) K. At temperatures below 650 K, the deviation is up to 10% down to 200 K and rises to 54% at 120 K.

4. COMPARISON WITH THEORETICAL DIELECTRONIC RECOMBINATION RATE COEFFICIENTS

Available theoretical C IV DR rate coefficients have been compiled recently by Savin (2000). For a critical assessment of the quality of the various calculations, the reader is referred to that work. Here, we confine ourselves to a brief comparison of our experimentally derived C IV DR plasma rate coefficient with the theoretical results, which in the following we divide into two categories.

1. The DR calculations by Burgess (1965), Shull & Steenberg (1982), Badnell (1989), and Chen (1991) have been carried out for high temperatures and do not reproduce the first local maximum of the experimental rate coefficient below 10,000 K (Fig. 5). At higher temperatures, the Burgess (1965) formula yields rate coefficients up to 50% larger and, at the other extreme, the result of Shull & Steenberg (1982) is 30% lower than our rate. Above 25,000 K the result of Badnell (1989) does not deviate more than 15% from ours. This deviation is within our experimental uncertainty. This is expected since Badnell also used the AUTOSTRUCTURE code for his DR calculations as we did for the high-\( n \) extrapolation of our measurement. The 15% deviation stems mainly from our scaling of the AUTOSTRUCTURE result and can be considered as the uncertainty of our experimentally derived rate coefficient at the temperature range where it is dominated by the AUTOSTRUCTURE extrapolation. The calculation of Chen (1991) gives rate coefficients only for temperatures higher than 10^5 K. They are up to 25% lower than our rate coefficient in that range.

2. Calculations that should be valid also at low temperatures have been published by McLaughlin & Hahn (1983), Nussbaumer & Storey (1983), Romanik (1988), Safronova, Kato, & Ohira (1997), and Mazzotta et al. (1998). They are shown in Figure 6 together with our result. The rate coefficient of Nussbaumer & Storey (1983) agrees with our result to within our systematic uncertainty of 15% only in the temperature range 4000–13,000 K; at lower temperatures it deviates up to 60%, and at higher temperatures it misses completely the second maximum of the rate coefficient that is due to DR via high Rydberg 1s^22pnl states. This is no surprise since the calculation by Nussbaumer & Storey (1983) has been restricted to low temperatures only. At temperatures ranging from 2000 to 20,000 K, the rate coefficient of Romanik (1988) deviates up to 60% from our
exponentially derived rate coefficient; at temperatures above 20,000 K, the deviation amounts to 37%. In this temperature range the result of Safronova et al. (1997) is almost a factor 2 lower than our derived rate. At temperatures below 20,000 K, it is up to a factor 1.5 larger. From 5000 to 5 \times 10^3 K, the result of McLaughlin & Hahn (1983) is within our error bar of 15%. The rate coefficient of Mazzotta et al. (1998) agrees reasonably well with our result only above 50,000 K. At temperatures below 2000 K, all available calculations deviate strongly from our experimentally derived rate coefficient. This is most probably due to the neglect of relativistic and many-body effects, which have been shown by Mannervik et al. (1998) to be essential for the correct description of the DR of even such a light ion as C IV. It should be mentioned again that at temperatures below \( \sim 10^3 \) K our result is completely independent of any theoretical model.

5. TOTAL RECOMBINATION RATE COEFFICIENT

Instead of separate calculations of RR and DR contributions, a unified treatment of both recombination processes can be considered. Such a treatment, which in principle also accounts for interference between DR and RR, has been presented by Nahar & Pradhan (1997). Experimentally, our measurement also yields the total recombination rate coefficient with the continuous RR contribution above 0.1 eV accurately represented by the fit described in § 3. Below that energy, all of the measured rate coefficient (Fig. 2) is assumed to be exclusively due to RR (see below). In principle, the total recombination rate coefficient can be derived by re-adding the continuous RR background to the composite DR spectrum, including the AUTOSTRUCTURE calculation-based extrapolation. However, at the low temperatures under consideration two experimental peculiarities require attention, namely, the much discussed (Gao et al. 1997; Gwinner et al. 2000) recombination rate enhancement at very low energies and the finite experimental resolution.

The measured enhancement of the recombination rate at very low energies is displayed in Figure 7. At energies below 3 meV, a strong enhancement of the experimental over the calculated rate coefficient sets in. At \( \tilde{E} = 0 \), this factor reaches a value of about 2.5. This effect has been found in electron-ion recombination measurements at different storage rings and has not been explained yet. Systematic studies of this effect have been carried out by Gao et al. (1997) and Gwinner et al. (2000). The recombination rate enhancement results most probably from the specific experimental arrangement at storage ring electron coolers where the electron beam is guided by a magnetic field.

Since we do not expect the enhancement to occur in an astrophysical environment, we subtract the excess rate coefficient (shaded area in Fig. 7) from our experimental data. To this end we have extrapolated the fitted RR background (see § 3) to lower energies by scaling an RR rate coefficient that has been calculated using a modified version of the semiclassical Bethe & Salpeter (1957) formula for the hydrogenic RR cross section; i.e.,

\[
\sigma^{RR}(\tilde{E}) = (2.10 \times 10^{-22} \text{ cm}^2) \times \sum_{n = n_{\text{min}}}^{n_{\text{max}}} t_q G_q(0) \frac{q^4 \tilde{r}^2}{n \tilde{E}(q^2 \tilde{r}^2 + n^2 \tilde{E})},
\]

with the Rydberg constant \( \tilde{r} \). For RR onto C IV, \( q = 3 \) and \( n_{\text{min}} = 2 \) are appropriate. As discussed in § 2, the maximum quantum number to be taken into account is determined by field ionization in our experimental setup. Since a rough estimate is sufficient for the present purpose, we take \( n_{\text{max}} = 20 \). The factors \( t_q \) account for partially filled shells. Here we use \( t_3 = \frac{7}{8} \) and \( t_q = 1 \) for \( n \geq 3 \). The Gaunt factors \( G_q(\tilde{E}) \) are small corrections that account for deviations of the semiclassical formula from the quantum mechanically correct hydrogenic result. Generally, the Gaunt factors are weakly energy dependent. We have taken energy-independent values calculated for \( \tilde{E} = 0 \) by Andersen & Bölkö (1990).

The calculated RR cross section has been convoluted by the experimental electron velocity distribution (see eqs. [4] and [9]), and the resulting rate coefficient has been multiplied by a constant factor of 1.638 such that it matches the fitted RR background at \( \tilde{E} = 0 \) eV. The finding that a factor different from unity has to be used in order to achieve the matching can be attributed to the fact that the hydrogenic treatment is not appropriate for RR, at least into low-\( n \) shells of the lithium-like C IV ion. The total recombination rate coefficient is now calculated as the sum of the extrapolated RR background (Fig. 7, solid line) and the DR rate coefficient derived in § 3 (see Table 1). This ensures that the excess rate due to the recombination rate enhancement does not contribute to the total recombination rate coefficient, which is displayed as the thick dash-dotted line in Figure 8.

At low plasma temperatures we also have to consider the influence of the finite experimental resolution (see eq. [5]) on our total plasma rate coefficient, which is well-defined only at temperatures \( T_e \gtrsim 120 \text{ K} \) as discussed before (§ 3). In order to quantify this influence, we have convoluted the theoretical RR cross section (eq. [13]) with \( n_{\text{max}} = 20 \) with our experimental energy distribution (see eqs. [4] and [9]) and derived from the resulting RR rate coefficient \( \sigma^{RR}(\tilde{E}) \) via equation (11) what we call here the “doubly convoluted” plasma RR rate coefficient. This is to be compared with the standard plasma RR rate, which is only “singly
where is the exponential

\[ \exp (-x) \]

account for RR into all \( n \) of high Rydberg states. In their calculation et al. Peüquignot et al. (1991) account for RR into all \( n \)-shells up to \( \infty \). For \( 4 \leq n < \infty \), they use the hydrogenic RR rate coefficient of Martin (1988). In order to compare our result to theoretical rate coefficients for finite values of \( n_{\text{max}} \), we first subtract this hydrogenic contribution from the rate coefficient of Peüquignot et al. (1991), thereby retaining their non-hydrogenic \( \text{C}^4 \) RR rate coefficient with \( n_{\text{max}} = 3 \). Onto this we have then added the hydrogenic RR rate coefficient calculated by using equation (14) with \( n_{\text{min}} = 4, n_{\text{max}} = 20 \) and \( n_{\text{min}} = 4, n_{\text{max}} = 40 \), yielding the thin dashed and dash-dotted curves in Figure 8, respectively. Extending the summation to \( n_{\text{max}} = 1000 \) yields—as expected—a curve that is indistinguishable from the result of Peüquignot et al. (1991) on the scale of Figure 8. At low temperatures our experimental rate coefficient approaches the theoretical curve that has been obtained with \( n_{\text{max}} = 40 \). This is roughly consistent with the prediction of the detailed model for field ionization in our experimental apparatus as discussed in § 2 and in the Appendix. From this observation we conclude that within our 15% systematic uncertainty the total \( \text{C}^4 \) recombination rate coefficient can be represented as the sum of the RR rate coefficient of Peüquignot et al. (1991) and our DR rate coefficient parameterized by equation (12) with the parameters listed in Table 1.

In Figure 9 this total \( \text{C}^4 \) recombination rate coefficient is compared to the results of the unified RR + DR calculation of Nahar & Pradhan (1997). At temperatures above 5000 K, both results agree with each other within our systematic uncertainty. At lower temperatures, however, pronounced differences occur. The theoretical result is up to a factor 3 larger than the experimental one. It should be noted that interference between RR and DR cannot be held responsible for the observed discrepancy. Interference is included in the unified calculations of Nahar & Pradhan (1997). And, it is also naturally included in the experimental result, which, as demonstrated above and shown in Figure 9, exhibits no measurable interference effect. We have a high degree of confidence in this conclusion as our experimental results include the strong 2p4\( \text{f} \) resonances where any interference effects are expected to be strongest. However, the lowest experimental \( \text{C}^4 \) DR resonance occurs at about 0.18 eV. Assuming a delta-like cross section and using the (somewhat too large) theory value of \( \sim 5 \times 10^{-19} \) cm\(^2\) eV for its strength (Mannervik et al. 1998), we calculate its contribution to the total rate coefficient by the use of equa-

![Figure 8](image_url)

**Figure 8.**—Experimental total \( \text{C}^4 \) recombination rate coefficients in a plasma corrected for the influence of the finite experimental resolution (thick solid line; systematic error \( \pm 15\% \)). The comparison with our pure DR rate coefficient (dotted line) shows that RR is noticeable up to \( \sim 30,000 \) K. The thick dash-dotted line is our total recombination rate coefficient uncorrected for the influence of the finite experimental resolution. The other lines are \( \text{C}^4 \) RR rate coefficients of Peüquignot et al. (1991; thin solid line) and corresponding RR rate coefficients (see text) for \( n_{\text{max}} = 20 \) (thin dashed line) and \( n_{\text{max}} = 40 \) (thin dash-dotted line).

![Figure 9](image_url)

**Figure 9.**—Total \( \text{C}^4 \) recombination rate coefficients in a plasma: this work (thick solid line; systematic error \( \pm 15\% \)) and theoretical unified calculation of Nahar & Pradhan (1997; dashed line). Our total recombination rate coefficient is obtained as the sum of the RR rate coefficient of Peüquignot et al. (1991; thin solid line) and our DR rate coefficient (dotted line; see eq. [12] and Table 1).
which field ionization may take place. The ion's survival probability is calculated from approximate field ionization rates inside the electron cooler up to the field ionization region (see Fig. 1). There the recombined ion spends a time during Bethe & Salpeter (1957). In equation (A1), which is valid for is measured from the center of the merging section Shore (1969), the magnitude of the plasma rate coefficient at low temperatures is sensitive to changes of such resonance positions on the milli-electron volt scale. As has convincingly been shown by Mannervik et al. (1998) and discussed above, the accurate prediction of resonance energies on this scale requires advanced and elaborate theoretical methods (Zong et al. 1997), which presently are not included in the more standard production codes for recombination rate coefficients of astrophysical interest. Consequently, at low plasma temperatures the DR and total rate coefficients resulting from such calculations should only be used with extreme care. As an example, we here have demonstrated that the low-temperature deviation of the theoretical prediction of Nahar & Pradhan (1997) from the experimental result is as large as a factor of about 3.

6. SUMMARY AND CONCLUSIONS

We have measured the $\Delta n = 0$ C IV DR rate coefficient by detecting with essentially full geometric efficiency the recombined ions produced following interaction in merged electron and ion beams. Generally, in such experiments the field ionization related to the charge analysis of the recombination products causes a large fraction of the DR resonance strength due to DR via high Rydberg states to remain unmeasured. For the present case of $\Delta n = 0$ DR of a low-charge ion, this undetected DR strength is substantial. In order to provide a remedy for this deficiency, we have performed AUTOSTRUCTURE calculations, carefully modeled the Rydberg-ion detection probabilities of our apparatus, scaled the theoretical results to our experimental low-energy DR rate coefficient, and, finally, used the high-energy part of the calculation as a substitution for the unmeasured DR strength. Using this extrapolated experimental DR spectrum, we have derived the C IV DR rate coefficient in a plasma (see eq. [12] and Table 1). Furthermore, the careful analysis of our experimental data leads to the conclusion that within our systematic uncertainty of 15% the total C IV recombination rate coefficient can be represented as the sum of the DR rate coefficient and the theoretical RR rate coefficient of Péquignot et al. (1991). This implies that interference between DR and RR is insignificant for the recombination of C IV. We have compared our rate coefficients with the available theoretical results. None of them agrees with our experimental rate coefficients over the full temperature range.

The field ionization of recombined ions in high Rydberg states, which is unavoidable in ion storage ring experiments, ultimately limits the capability of providing meaningful, pure experimental plasma DR rate coefficients. This is especially true for low-charge-state ions where much DR strength is concentrated in high-$n$ resonances. The DR resonance strength drops much faster with increasing $n$ for high-charge-state ions so that a smaller fraction of the total DR strength is accumulated in high Rydberg states. Consequently, for higher charge states storage ring experiments can provide reliable DR plasma rate coefficients with much more limited need of extrapolation.

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APPENDIX A

MODEL CALCULATION OF DETECTION PROBABILITIES FOR HIGH RYDBERG STATES

In the detailed model of the field ionization properties of our apparatus, the survival fractions are determined individually for each $nl$-state populated by DR, in combination with the flight times $t_F$ to the field ionization zones and hydrogenic decay rates. Assuming a constant recombination probability across the length $L$ of the merging section, the probability that a state characterized by quantum numbers $n$ and $l$ has decayed upon reaching the field ionization zone is given as

$$P_d(nl, t_F, t_p) = 1 - \frac{\tau(nl)}{t_L} \left[ \exp \left( -\frac{t_p - t_F/2}{\tau(nl)} \right) - \exp \left( -\frac{t_p + t_F/2}{\tau(nl)} \right) \right].$$  (A1)

The flight time through the electron cooler is $t_L = 87$ ns in the present case, and the radiative lifetime of the $nl$-state is calculated as $\tau(nl) = \left[ \sum_{n'<n, l'=l} \gamma_{nl,n'l} \right]^{-1}$ from hydrogenic radiative dipole transition rates $\gamma_{nl,n'l}$ following Bethe & Salpeter (1957). In equation (A1), which is valid for $t_F \geq t_p/2$, $t_F$ is measured from the center of the merging section inside the electron cooler up to the field ionization region (see Fig. 1). There the recombined ion spends a time $\Delta t_F$ during which field ionization may take place. The ion's survival probability is calculated from approximate field ionization rates $A_F$ derived by Damburg & Kolosov (1979) to be

$$P_d(nl, \Delta t_F, F) = \frac{1}{2l + 1} \sum_{m=-l}^{l} \sum_{n_1=0}^{\left| n - m \right| - 1} \left( C_{n_1,m}^{n,l} \right)^2 \exp \left[ -\Delta t_F A_F(n_1, n_2, m) \right],$$  (A2)

where the expansion coefficients are the Clebsch-Gordan coefficients

$$C_{n_1,m}^{n,l} = \left\langle \begin{array}{c} (n-1)/2 \\ m + n_1 - n_2 \\ (m-n_1 + n_2)/2 \end{array} ; \begin{array}{c} (n-1)/2 \\ m \\ l \end{array} \right\rangle.$$  (A3)
for the transformation to Stark states characterized by the parabolic quantum numbers and $m$, which satisfy the relation $n = n_1 + n_2 + |m| + 1$ (Gallagher 1994). Finally, the detection probability is calculated from the decay and survival probabilities as

$$\gamma_n(t_L, t_F, \Delta t_F, F) = \left[ 1 - P_d(nl, t_L, t_F) \right] P_s(nl, \Delta t_F, F) + \sum_{n'<n, |l|=l'=\pm 1} b(nl \rightarrow n'l')$$

$$\times \left\{ [P_d(nl, t_L, t_F) - P_d(nl, n'l', t_L, t_F)] P_s(n'l', \Delta t_F, F) + [P_{sc}(nl, n'l', \ldots, t_L, t_F, \Delta t_F, F)] \right\},$$

(A4)

where branching ratios for a dipole transition from state $nl$ to state $n'l'$ are defined as $b(nl \rightarrow n'l') = \tau(nl)f(nl \rightarrow n'l')$. The quantity $P_{sc}(nl, n'l', \ldots, t_L, t_F, \Delta t_F, F)$ accounts for cascading, i.e., stepwise de-excitation by more than one transition. It is calculated recursively. In equations (A4) and (A5),

$$P_d(n_1 l_1, \ldots, n_N l_N, t_L, t_F) = \sum_{k=1}^{N} \sum_{n'=n_l}^{n_1} \frac{[\tau(n_l l_k)]^{n_l - 1}}{[\tau(n_k l_k) - \tau(n_l l_k)]} P_d(n_k l_k, t_L, t_F)$$

(A6)

is the probability for a decay along a given sequence $n_1 l_1, \ldots, n_N l_N$ of $N$ excited hydrogenic states (Schippers 1995). The one-step decay probabilities $P_d(n_k l_k, t_L, t_F)$ are evaluated according to equation (A1).

Due to the high $n$-values to be considered (up to $n \approx 100$), the explicit calculation of cascade contributions is very laborious and time consuming. In view of the fact that rates are highest for transitions to the lowest available state and decrease with $\sim n^{-3}$ (Bethe & Salpeter 1957), it can be anticipated that except for high angular momentum states, which are only sparsely populated by DR, cascading plays only a minor role. We have verified this by comparing calculations without and with cascades via one and two intermediate states. The resulting rate coefficients are indistinguishable on the scale of Figure 3. Therefore, we conclude that it is safe to neglect contributions from cascades with more intermediate states.

In the calculation of detection probabilities, we include states up to $n = 100$, which is far beyond the estimated hard cutoff at $n_F = 19$ (see §2). Furthermore, we consider all field ionization regions that have been mentioned previously: the toroidal section at the exit of the electron cooler, two correction magnets, and the charge-analyzing dipole magnet (Fig. 1). The corresponding detection probabilities $\gamma_n^{(0)}$, $\gamma_n^{(1)}$, $\gamma_n^{(2)}$ and $\gamma_n^{(3)}$, respectively, are calculated individually, and the overall detection probability is finally given as the product of the four individual ones, i.e.,

$$\gamma_n = \gamma_n^{(0)}(t_L, t_F, \Delta t_F, F^{(0)}) \times \gamma_n^{(1)}(0, t_F^{(1)} - t_F^{(0)}, \Delta t_F^{(1)}, F^{(1)})$$

$$\times \gamma_n^{(2)}(0, t_F^{(2)} - t_F^{(1)}, \Delta t_F^{(2)}, F^{(2)}) \times \gamma_n^{(3)}(0, t_F^{(3)} - t_F^{(2)}, \Delta t_F^{(3)}, F^{(3)}),$$

(A7)

with the numerical values for $t_F$, $\Delta t_F$, and $F$ taken from Table 2. It should be noted that for the calculation of $\gamma_n^{(1)}$, $\gamma_n^{(2)}$, and $\gamma_n^{(3)}$, we used the $t_L \rightarrow 0$ limit of equation (A1), namely, $P_d(nl, 0, t_F) = 1 - \exp \left\{ -t_F/\tau(nl) \right\}$. A contour plot of the calculated detection probabilities is shown in Figure 4. The detection probabilities are 100% for $nl$-states up to $n = 19$. At higher $n$, only low-$l$ states have a high detection probability owing to their short lifetimes. Higher $l$ states have increasingly higher lifetimes and are more and more effectively field ionized. Consequently, for $n > 19$ the detection probability drops very fast to zero with increasing $l$. Above $n \approx 40$, low-$l$ states are also cutoff due to field ionization in the toroid. According to equation (6) the hard cutoff number there is $n_F = 45$. Due to the short time of only 55 ns needed for traveling from the cooler center to the toroid, almost no radiative de-excitation of higher Rydberg states is possible.

### Table 2

**Field Ionization Regions in the Flight Path of Recombined Ions**

| Region                  | $F$  | $\Delta t_F$ | $t_F$ | $n_F$ |
|-------------------------|------|--------------|-------|-------|
| Toroid                  | 4    | 29           | 55    | 45    |
| First correction magnet | 6    | 20           | 103   | 40    |
| Second correction magnet| 12   | 10           | 133   | 34    |
| Dipole magnet           | 106  | 49           | 275   | 19    |

**Notes.**—Units are kV cm$^{-1}$ for $F$ and nanoseconds for $\Delta t_F$ and $t_F$. The term $n_F$ has been calculated from eq. (6).
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