[10]For a semiclassical "explanation" of the linewidths that give rise to this height ratio in the Mollow case, see C. Cohen-Tannoudji, in *Laser Spectroscopy*, edited by S. Haroche *et al.* (Springer, Berlin, 1975), p. 336.

[11]Results in this case have already been reported by the author in Proceedings of the Symposium on Resonant Light Scattering, Cambrdige, Massachusetts, April 1976 (unpublished). The subsequent incorporation of phase-fluctuation effects into the resonance-fluorescence theory has also been achieved by G. S. Agarwal (private communication). See also remarks by Cohen-Tannoudji, Ref. 10.

[12]K. Wódkiewicz and J. H. Eberly, in preparation, and Ref. 6 (to be published).

Enhancement of Dielectronic Recombination by Plasma Electric Microfields*

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An investigation has been made of the l-redistribution effect of quasistatic ion electric microfields on dielectronic recombination. Using the linear-Stark-effect approximation, a hydrogen ion density of 10^{14} cm^{-3} is found to produce a threefold enhancement in the rate for the dielectronic recombination process $Fe^{+23}(2s) + e^- \rightarrow Fe^{+22}(2p, nl) \rightarrow Fe^{+22}(2s, nl) + \hbar\omega$.

In this Letter we report on the application of techniques from Stark-broadening theory[1] to calculate the dielectronic recombination rates for multiply charged Fe ions under the influence of plasma electric microfields. Results are presented for a dielectronic recombination process where the effects of the surrounding charged particles have been found to be important at surprisingly low densities.

The conventional interpretation of certain spectral line intensities may be distorted, or even invalidated, by the combined effect of dielectronic-recombination satellites[2] whose wavelengths are indistinguishable from that of the associated resonance line from the recombining ion. In addition, dielectronic recombination has been shown[3] to be the dominant recombination process for nonhydrogenic impurity ions in low-density high-temperature plasmas, such as the solar corona and the discharges produced in controlled thermonuclear fusion experiments.

Dielectronic recombination is the result of a radiationless capture into a doubly excited state

$$X^{+Z}(i) + e^-(\epsilon_i) \rightarrow X^{+(Z-1)}(j, nl), \tag{1}$$

followed by a stabilizing radiative transition involving de-excitation of the recombining ion core

$$X^{+(Z-1)}(j, nl) \rightarrow X^{+(Z-1)}(k, nl) + \hbar\omega. \tag{2}$$

At low densities, where the final excited ions cascade to their ground states in times that are short compared with the electron-ion collision time, the overall recombination rate simply equals the total rate for all stabilizing radiative transitions.

Burgess and Summers[4] have pointed out that with increasing density the overall recombination coefficient defined by Bates, Kingston, and McWhirter[5] is reduced by the effects of collisional ionization on the highly excited nl states, which can play the most important role in the dielectronic recombination process. Burgess and Summers[4] were also the first to suggest that the dielectronic recombination rate could be enhanced by the effects of collisionally induced angular-momentum redistribution of the doubly excited states.

The spectral intensity arising from the stabilizing radiative transitions may be calculated by employing the techniques of modern Stark-broadening theory,[1] which must be applied in its most general form in order to describe the effects of perturbing electrons and ions on the nearly dengerate l sublevels of the upper and lower states. The action of the ions is customarily treated in the quasistatic and long-range dipole (Stark effect) approximations.[6] The effects of electron collisions may be taken into account by applying the generalized impact approximation[7-10] to the quasistatic Stark components before performing the average over the ion electric microfield distribution.

In the present investigation, account is taken only of mixing of the nearly degenerate l sublevels due to the action of the quasistatic ion electric microfields. We defer to a future investigation the inclusion of electron collisions, which will become important with increasing density.

The spectral intensity emitted by the set of stabilizing radiative transitions $(j, n) \to (k, n)$ is obtained as a superposition of the intensities arising from the individual quasistatic Stark components, which are weighted by the probability that the fluctuating electric field produces the appropriate displacement of the photon frequency. In the customary approximation[3] of ignoring the effect of the $j \to k$ transition on the state of the outer electron, there is no displacement in the emitted photon frequency from the $j \to k$ resonance line center. This effect might have to be taken into account to explain the marked blue (plasma-polarization) shift which has recently been observed[11] in a He^+ Lyman-α emission line and which could be in part attributable to a blend of dielectronic-recombination satellites that have been Stark-shifted from their normal position on the long-wavelength side of the resonance line.[1]

The most important effect of the electric fields on dielectronic recombination is revealed when the doubly excited level populations are determined in the Stark representation. In the linear-Stark-effect approximation, the outer-electron states are specified by $n\lambda m$, where λ is the electric quantum number. The doubly excited level populations are determined by the radiationless capture rate $C(i, \epsilon_i \to j, n\lambda m)$ and by the total decay rate $A(j, n\lambda m)$. The total decay rate is the sum of the autoionization rates $A_a(j, n\lambda m \to k, \epsilon_k)$ into all open continuum channels (k, ϵ_k) and of the radiative decay rates $A_r(j \to k)$ into all lower states k. Integration over the Lorentzian profiles associated with the total decay widths, which are assumed to remain small in comparison with the level separations, yields the total emission rate in the form $N(i)N_e\alpha_d(j, n \to k, n)$, where $N(i)$ is the initial ion density and N_e is the electron density. The recombination coefficient $\alpha_d(j, n \to k, n)$ is obtained in the form

$$\alpha_d(j, n \to k, n) = 2^3 a_0^{\,3}\pi^{3/2}(E_H/k_B T_e)^{3/2}\frac{g(j)}{g(i)}A_r(j \to k)\sum_{\lambda m}\frac{A_a(j, n\lambda m \to i, \epsilon_i)}{A(j, n\lambda m)}\exp\left[\frac{E(i) - E(j) - E(n\lambda m)}{k_B T_e}\right], \tag{3}$$

where $g(i)$ and $g(j)$ are the statistical weights associated with the energy levels $E(i)$ and $E(j)$ and $E_H = 13.6$ eV. The phenomenon of field-induced autoionization, which has been treated previously[12] in the quadratic-Stark-effect approximation, is described in the linear-Stark-effect approximation by the transformation

$$A_a(j, n\lambda m \to i, \epsilon_i) = \sum_{l=|m|}^{n-1} |\langle m\lambda m | nlm\rangle|^2 A_a(j, nl \to i, \epsilon_i). \tag{4}$$

The field-free autoionization rates $A_a(j, nl \to i, \epsilon_i)$ have been found[13] to be negligible for $l > 15$. From the relationship between the corresponding capture and autoionization rates, it follows that the electric fields can induce dielectronic recombination through the normally inaccessible higher-angular-momentum states, which have larger statistical weights.

The validity of the linear-Stark-effect approximation can be insured by restricting the use of Eq. (3), whose field-strength dependence may be neglected, to a range of outer-electron principal quantum numbers which is to be determined for each value of the mean field strength. The lower limit on n may be taken to be the value for which twice the m-averaged quadratic Stark level shift,[12] obtained for a representative l value, becomes equal to the l sublevel separation. For the high l values of interest, the sublevel separation is determined by long-range interactions with the multipole moments of the ionic core and by fine-structure splitting. The upper limit on n can be obtained from the Inglis-Teller criterion[14] that twice the largest linear Stark level shift must not exceed the separation between levels with adjacent n values. Above the Inglis-Teller limit, account must be taken of higher-multipole radiator-perturber interactions, mixing of different n levels, and field-induced ionization due to tunneling.[15]

In a thermal plasma, the mean field strength F produced by a given perturber species p may be related to the density N_p by means of the Holtsmark relationship,[1] $F = 8.8 Z_p e N_p^{\,2/3}$. Only the hydrogen ion contribution to F has been taken into account in the present calculation, but it should be noted that the Fe impurity ions may provide an important contribution in some controlled fusion experiments.

Equations (3) and (4) have been employed to calculate the rate coefficients for dielectronic recombination of Fe^{+23} ions through their $2p \to 2s$ transition (with $j = 2p$ and $i = k = 2s$). Autoionization rates were obtained from distorted-wave results for the threshold partial-wave $2s$-$2p$ collisional excitation cross sections.[13] The electron temperature was chosen to be the value 3×10^7 °K, where the Fe^{+23} ion is found[13] to have its maxi-

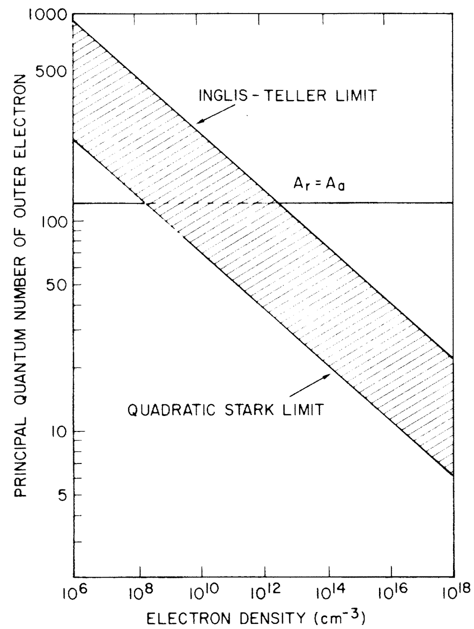

FIG. 1. Validity n region of the linear-Stark-effect approximation for the $2p \rightarrow 2s$ dielectronic recombination of Fe^{+23} ions.

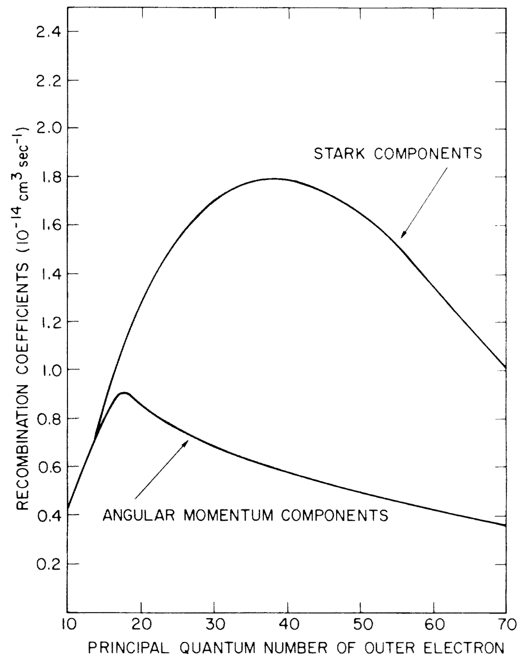

FIG. 2. Comparison between the Fe^{+23} $2p \rightarrow 2s$ dielectronic recombination rate coefficients obtained by summation over linear Stark components and angular-momentum components.

mum abundance in an isothermal steady-state plasma. The validity n region of the linear-Stark-effect approximation is shown in Fig. 1 for a range of electron (or hydrogen ion) densities. The fact that the radiative decay rate exceeds the autoionization rate only for $n > 123$ indicates the importance of large n values in the case of a $\Delta n = 0$ transition of the recombining ion. For this highly charged ion, the additional $3p \rightarrow 2s$ dielectronic recombination process is found[13] to occur primarily through the $n = 3$ transition and, consequently, would be influenced by the electric fields only at densities exceeding 10^{18} cm^{-3}.

The results for $\alpha_d(2p, n \rightarrow 2s, n)$ obtained using the linear Stark representation are compared in Fig. 2 with the corresponding field-free results, which were obtained by employing the usual l representation. The use of the linear Stark representation is found to produce a shift in the intensity maximum from $n = 17$ to $n = 38$ together with a threefold enhancement.

The total rate of the $2p \rightarrow 2s$, radiatively stabilized transitions can be estimated conservatively by using the field-free results below the quadratic Stark limit and also the linear-Stark-effect results between the quadratic Stark and the Inglis-Teller limits. For an electron density of 10^{14} cm^{-3} and an electron temperature of 3×10^7 °K, the total rate coefficient is found to be

$$\alpha_d(2p \rightarrow 2s) = 0.84 \times 10^{-12} \text{ cm}^3 \text{ sec}^{-1} \quad \text{(Stark components)}. \tag{5}$$

This is to be compared with the field-free result

$$\alpha_d(2p \rightarrow 2s) = 0.36 \times 10^{-12} \text{ cm}^3 \text{ sec}^{-1} \quad \text{(angular-momentum components)}, \tag{6}$$

which is close to the value 0.47×10^{-12} cm^3 sec^{-1} obtained[13] from the zero-density Burgess formula,[16] indicating that the effect of the Inglis-Teller cutoff is not important. The total rate coefficient for the just as important $3p \rightarrow 2s$ process is found[13] to be 0.96×10^{-12} cm^3 sec^{-1}, which is reduced by a factor of 10 from the value predicted by the Burgess formula[16] because of the inclusion of autoionization into the $3s$ excited state of the recombining Fe^{+23} ion.[13]

It appears desirable to calculate the overall recombination rate coefficient defined by Bates, Kingston, and McWhirter[5] by combining the procedure of Burgess and Summers,[4] which describes the effects of collisions on the populations of the final states in the stabilizing radiative transitions, with a treatment of the influence of the plasma on the doubly excited level populations.

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Formation of Potential Double Layers in Plasmas*

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Large localized potential double layers, $e\varphi >> kT_e$, are observed to develop from a solitary E-field structure in a plasma traversed by electron drifts, $v_d \simeq v_e$. Strong electron and ion accelerations occur across the layers. At large electron drift velocities ($v_d > 3v_e$), layers are unstable to large ion fluctuations with $\delta n/n$ approaching unity.

The existence of intense localized dc electric fields, potential double layers, has been the subject of many investigations in connection with intense beam injections[1] and in the ionospheric formation of auroral substorms.[2] We wish to present experimental evidence on the formation and stability of such potential layers in a bounded low-density plasma ($n \sim 10^8$ cm^{-3}) traversed by electron drifts, $v_d \approx v_e$ [$v_e = (kT_e/m)^{1/2} \sim 0.8 \times 10^8$ cm/sec]. Coherent potential perturbations initially excited by electron-ion drift instability[3,4] are observed to transform into stationary potential jumps[5] which exist self-consistently with the interpenetrating plasma streams and reflected ions and electrons.

The experiments were performed in the University of California at Los Angeles double-plasma beam system[6,7] which consists of two electrically isolated stainless-steel cylinders separated by two closely spaced tungsten grids (with 200×200 fine mesh). Plasmas are produced by bombardments of neutral gas (Argon, $P_0 \sim 10^{-4}$ Torr) with electrons from a set of hot filaments in the source chamber. The plasma potentials can be varied by applying an appropriate voltage bias between the two chambers to produce plasma flows. Three distinct regimes of operations can be accomplished, as shown in Fig. 1. *Regime I:* Slow electron-ion drifts [$0.3v_e > v_d > v_s$, where $v_e = (kT_e/m)^{1/2}$ and $v_s = (kT_e/M)^{1/2}$ are the electron thermal and ion acoustic speeds, respectively]. Plasma fluctuations identified as the current-driven ion acoustic instabilities were observed to grow spatially and become saturated at $\delta n/n \leq 10\%$ by an electron trapping process. Small amplitude waves were found[6] to develop into large-amplitude

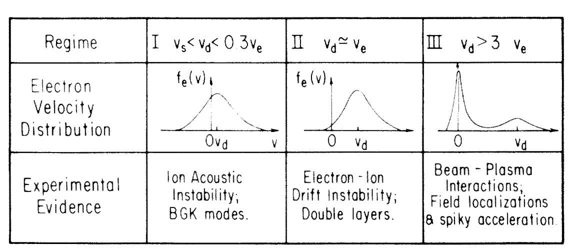

FIG. 1. The three regimes of plasma interactions observed with use of the double-plasma beam system.