Collision quenching of Rydberg atomic levels and electron-ion recombination in a noble buffer gas

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Three-particle recombination of electrons with atomic Xe^+ ions in a plasma of a noble-gas mixture He/Xe ($N_{Xe} \ll N_{He}$) is investigated experimentally and theoretically. At ionization degrees $\alpha \sim 10^{-8}$–10^{-5} and electron temperatures $T_e \gtrsim 700$–1000 K the recombination rates decrease much more slowly with increase of T_e, and they are correspondingly much larger than the rates calculated using the diffusion mechanism of three-particle recombination on free electrons of the plasma. It is also established that in this range of plasma parameters the recombination flux is proportional to the He buffer-gas concentration. This points to a substantial role of neutral particles in electron-ion recombination. The traditional Fermi mechanism of energy quenching of a highly excited "recombining" electron by buffer-gas atoms does not explain the observed anomalous behavior of the recombination coefficients. A new quenching mechanism is proposed and investigated, brought about by nonadiabatic energy exchange between highly excited electrons and the internal electrons of a quasimolecular heteronuclear noble-gas ion produced upon scattering of a perturbing neutral particle (He) by the ion core (Xe^+) of a Rydberg atom ($Xe(n)$). A recombination model is constructed, in which account is taken of the diffusion over the energy levels of the ($Xe(n)$) atom on collision with the plasma free electrons, as well as of the effective quenching of the Rydberg-electron energy by the buffer gas (He) atoms. The experimental results of the paper are explained quantitatively on the basis of this model.

1. INTRODUCTION

It is known that in a low-temperature plasma three-particle recombination of electrons (e^-) with atomic ions (X^+) proceeds as a rule via formation of $X(nl)$ atoms in Rydberg nl states, and has the character of diffusion over highly excited energy levels $nl \to nl'$ in collisions with free electrons e^- or with neutral particles of a buffer gas Y (or X itself):

$$X^+ + e^- + e^- \to X(nl) + e^-, \quad (1a)$$

$$X(nl) + e^- \to X(n'l') + e^-, \quad (1b)$$

$$X^+ + e^- + Y \to X(nl) + Y, \quad (2a)$$

$$X(nl) + Y \to X(n'l') + Y. \quad (2b)$$

The rates of electron-ion recombination via collisions (1) and (2) with the plasma electrons[1] and with buffer-gas atoms[2] were calculated in the diffusion approximation by Gurevich and Pitaevskiĭ.[1,2] In Ref. 2 (and in most succeeding studies of recombination in collisions with neutral particles) the inelastic collisional transitions between the highly excited levels ($nl \to nl'$) of the $X(nl)$ atom in the course of the recombination were considered in the framework of the Fermi mechanism due to scattering of the perturbing atom by the weakly bound electron e^-. The role of the ionic core X^+ in this mechanism reduces only to production of a specified (for the considered nl level of the atom $X(nl)$ distribution of the momenta of the weakly bound electron in a Coulomb field. Specific calculations[1-3] show that at sufficiently high degrees α of plasma ionization the three-particle electron-ion recombination takes place in collisions (1) with free electrons, while collisions (2) with the buffer-gas atoms Y, via the Fermi mechanism, are substantial only at the very low values $\alpha \lesssim 10^{-8}$.

We have investigated experimentally and theoretically the recombination of electrons with atomic X^+ ions in a plasma of mixed inert gases X/Y ($N_X \ll N_Y$). In this process collisions with the neutral Y particles play a decisive role even at rather high degrees of ionization $\alpha = N_e/N_Y = 10^{-8}$–$10^{-4}$ (which are typical of a plasma produced by an electron beam or by a pulse discharge). The actual experiments were performed for afterglow plasma of a pulse discharge in a helium-xenon mixture ($N_{He} = 2.6 \cdot 10^{17}$–$4.4 \cdot 10^{18}$ cm^{-3}, $N_{Xe} = 10^{14}$–10^{16} cm^{-3}, $T = 300$–600 K). Under these conditions the dissipative recombination of the molecular ions Xe_2^+ can be neglected, since $[Xe_2^+] \ll [Xe^+] \approx N_e$. The method proposed in the present paper yielded the dependences of the coefficients $\beta(T_e, N_e)$ [cm$^6 \cdot$s^{-1}] of three-particle recombination of electrons with Xe^+ atomic ions and of the emission intensities J_λ (T_e, N_e) of a number of Xe-atom spectral lines on the temperature and on the electron density in the ranges $T_e = 400$–2500 K and $N_e = 2 \cdot 10^{10}$–$3 \cdot 10^{12}$ cm^{-3}. A feature of the obtained dependences (see Fig. 4) is a steep decrease (close to $T_e^{-9/2}$) in the region of low temperatures $T_e \lesssim 800$ K, and a substantially slower decrease and anomalously high recombination

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rates at $T_e \gtrsim 800$ K. It was also established that the recombination flux $\Gamma[\mathrm{cm}^{-3}\cdot\mathrm{s}^{-1}] = \beta N_e^3$ depends substantially on the degree of plasma ionization and on the helium concentration.

This anomalous character of the recombination cannot be explained by using the known models of electron recombination with atomic ions in three-particle collisions with electrons or with atoms of the buffer gas He. Electron recombination with Xe^+ ions in collisions with atoms of Xe itself (Ref. 4), via dipole transitions between symmetric and antisymmetric terms ($\Sigma_g \to \Sigma_u$) of the quasimolecular ion Xe_2^+, is also ineffective here because of the low xenon concentration. To explain the experimental results we have therefore constructed in the present paper a recombination model in which the capture of the recombining electron e^- by the inert-gas atomic ion X^+ and its relaxation over the highly excited levels of the $X(nl)$ atom is initially mainly by diffusion through collisions with electrons (just as in the model of Ref. 1) and subsequently, starting with certain levels $|\varepsilon_{nl}| \gtrsim \varepsilon_Y$, predominantly via effective quenching by atoms of the buffer gas Y. The traditional Fermi quenching mechanism, due to scattering of the perturbing atom Y by a weakly bound electron e^- of a highly excited atom $X(nl)$, in the principal quantum-number (n) region of importance for recombination, is ineffective in this case.

We propose here therefore a new mechanism of quenching highly excited levels, realized by collisions $X(nl) + Y$ of different noble-gas atoms and due to scattering of the perturbing particle $Y(^1S_0)$ by the ionic core $X^+[n_0p^5(^2P_{3/2})]$. This mechanism is the result of quadrupole and short-range parts of the Coulomb interaction between the external electron e^- and the internal electrons of the quasimolecular ion YX^+ produced when the perturbing atom Y passes near the ionic core X^+. Collision de-excitation $nl \to n'l'$ ($\varepsilon_{nl} > \varepsilon_{n'l'}$) of a Rydberg electron e^- is then the result of nonadiabatic transitions $X|j, \Omega\rangle \to A_1|j, \Omega'\rangle$ in the quasimolecular heteronuclear ion YX^+ of the inert gas, between its lower split terms $U_{j\Omega}(R)$ and $U_{j\Omega'}(R)$ (see Fig. 1). These terms correlate with the ground states of the isolated atom and ion of the noble gases $X^+(n_0p^5, {}^2P_{3/2}) + Y(^1S_0)$, but have different values of the components $\Omega = 1/2$ and $\Omega' = 3/2$ of the total angular momentum $\mathbf{j} = \mathbf{L} + \hat{\mathbf{S}}(j = 3/2)$ of the inner electron shell along the internuclear axis $\mathbf{R}$. Note that the term $A_2|j' = 1/2, \Omega' = 1/2$, and all the more other excited terms of the YX^+ ion do not participate in the investigated mechanism at thermal energies $T \sim 0.03$–0.05 eV ($T \ll \Delta_{jj'}$) of collision with the Y atoms, in the case of heavy noble-gas ions X^+ having a large spin-orbit splitting $\Delta_{jj'}$ of the states $^2P_{3/2}$ and $^2P_{1/2}$ of the electronic n_0p^5 shell ($n_0 = 3$, 4, and 5 respectively for Ar^+, Kr^+, and Xe^+).

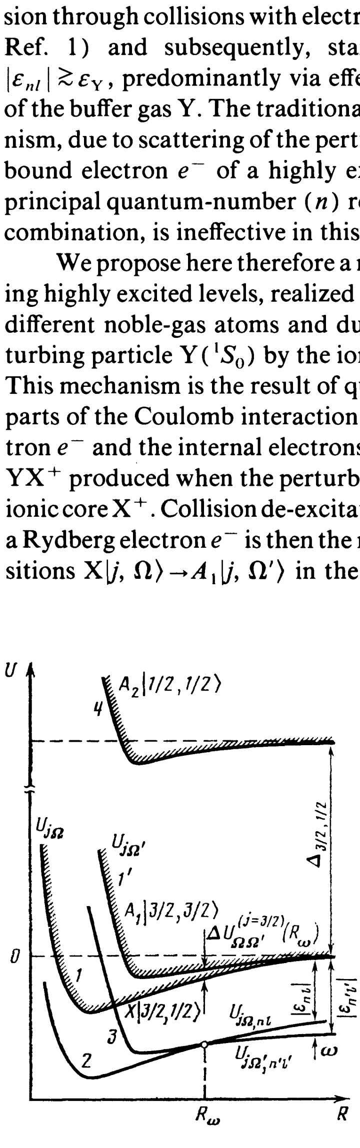

FIG. 1. Typical form of lower electronic terms of heteronuclear noble-gas ions YX^+ and of the corresponding Rydberg terms of the quasimolecule $YX^+ + e^-$, with the system $HeXe^+ + e^-$ as an example. Curves 1, 1', and 4—respectively the ground, first, and second excited terms of the $HeXe^+$ ion, which correlate with the states $Xe^+[5p^5(^2P_{3/2})] + He(^1S_0)$ (curves 1 and 1') and $Xe^+[5p^5(^2P_{1/2})] + He(^1S_0)$ (curve 4) of the isolated atom and ion. Curves 2 and 3—Rydberg terms that correlate with the states of the correlated atoms in the initial $Xe[5p^5(^2P_{3/2}), nl] + He(^1S_0)$ (curve 2) and final $Xe[5p^5(^2P_{3/2}), n'l'] + He(^1S_0)$ (curve 3) states.

Within the scope of the constructed recombination model, which takes into account the collisions with free plasma electrons and the effective quenching of the Rydberg atomic levels by the buffer-gas electrons, we present a quantitative explanation of the experimentally observed (see Sec. 4) effect of the strong increase of the rate of electron-ion recombination in the He/Xe mixture plasma compared to ordinary three-particle recombination on electrons.

2. INELASTIC TRANSITIONS BETWEEN RYDBERG LEVELS IN COLLISIONS OF UNLIKE NOBLE-GAS ATOMS

2.1. Formulation of problem

In the mechanism investigated, $nl \to nl'$ transitions between Rydberg levels of an $X(nl)$ atom by passage of a perturbing atom Y near the ionic core X^+ are produced in the quasimolecular ion YX^+ at interbarrier distances R much smaller than the orbit radius $\langle r_e \rangle \sim 2n^2$ of the highly excited electron e^-. It is therefore convenient to express the total Hamiltonian H of the entire $YX^+ + e^-$ system in the form (the units $e = m_e = \hbar = 1$ are used in this paper)

$$\begin{gathered} H = T_{\mathbf{R}} + H_{YX^+}(\mathbf{r}_\varkappa, \mathbf{R}) + H_e(\mathbf{r}_e) + V, \\ H_e|\psi_{nlm}(\mathbf{r}_e)\rangle = \varepsilon_{nl}|\psi_{nlm}(\mathbf{r}_e)\rangle, \\ H_{YX^+}|\varphi_{j\Omega}(\mathbf{r}_\varkappa, \mathbf{R})\rangle = U_{j\Omega}(R)|\varphi_{j\Omega}(\mathbf{r}_\varkappa, \mathbf{R})\rangle. \end{gathered} \tag{3}$$

Here $T_{\mathbf{R}} = -\Delta_{\mathbf{R}}/2\mu$ is the operator of the kinetic-energy of the nuclei (μ is their reduced mass); H_{YX^+} is the Hamiltonian of the electron shell of the heteronuclear ion YX^+; $\varphi_{j\Omega}(\mathbf{r}_\varkappa, \mathbf{R})$ are its proper wave functions corresponding to ths considered adiabatic terms $U_{j\Omega}(R)$ and $U_{j\Omega'}(R)$ with $j = 3/2$, $\Omega = 1/2$ and $j = 3/2$, $\Omega' = 3/2$; $H_e = -\Delta_{\mathbf{r}e}/2 - 1/r_e$ is the Hamiltonian of the Rydberg electron e^- in the Coulomb field of the ionic core X^+, and $\psi_{nlm}(\mathbf{r}_e)$ are its proper wave functions (n, l, and m are the principal, orbital, and magnetic quantum numbers) with energy $\varepsilon_{nl} = -1/2n_{\text{eff}}^2 n_{\text{eff}} = n - \delta_l, \delta_l$ is the quantum defect, and $V = \Sigma_\varkappa |\mathbf{r}_e - \mathbf{r}_\varkappa|^{-1}$ is the operator of the interaction of the external ($\mathbf{r}_e$) and internal ($\mathbf{r}_\varkappa$) electrons of the $YX^+ + e^-$ system.

This interaction is responsible for the investigated mechanism of the process (2b) of the collisional quenching $nl \to n'l'$ of the Rydberg electron e^-, accompanied by the nonadiabatic transition $X|j = 3/2, \Omega = 1/2\rangle \to A_1|j = 3/2, \Omega' = 3/2\rangle$ in the electron shell of the quasimolecular ion YX^+ between its lower split terms $U_{j\Omega}(R)$ and $U_{j\Omega'}(R)$. This nonadiabatic energy exchange between the external and internal electrons takes place in the immediate vicinity of the intersection point R_ω of the Rydberg terms $U_{j\Omega,nl}(R) = U_{j\Omega}(R) = +\varepsilon_{nl}$ and $U_{j\Omega',n',l'}(R) = U_{j\Omega'}(R) + \varepsilon_{n'l'}$ of the $YX^+ + e^-$ quasimolecule. In other words, the collision quenching of the highly excited atom ($nl \to n'l'$) takes place in a region of internuclear distances near the point R_ω, for which the splitting

$$\Delta U^{(j)}_{\Omega\Omega'}(R_\omega) = U_{j\Omega'}(R_\omega) - U_{j\Omega}(R_\omega)$$

of the electron terms of the heteronuclear quasimolecuar ion YX^+ (produced in the course of the collision between the perturbing atom Y and the ionic core X^+ becomes equal to the frequency $\omega = \varepsilon_{nl} - \varepsilon_{n'l'}$ of the transition of the Rydberg electron e^- (see Fig. 1).

Within the framework of the stationary perturbation theory with respect to Coulomb interaction of the external and internal electrons of the $YX^+ + e^-$ system we have for the cross section of theprocess (2b) of inelastic collision transitions $nl \to nl'$ between the Rydberg levels of the atom, in the semiclassical approximation,

$$\sigma^{\Omega\Omega'}_{nl,n'l'} = \frac{4\pi^2\mu^2 g_{j\Omega}}{(2l+1) g_{tot} q q'} \sum_{mm'} \int_0^\infty 2\pi\rho \, d\rho \, |\langle \chi_{j\Omega,q\rho} | \langle \varphi_{j\Omega} | \times \langle \psi_{nlm} | V | \psi_{n'l'm'} \rangle | \varphi_{j\Omega'} \rangle | \chi_{j\Omega',q'\rho} \rangle |^2. \quad (4)$$

Here $g_{j\Omega} = 2$ and $g_{j\Omega'} = 2$ are the statistical weights of the ground $X|j = 3/2, \Omega = 1/2\rangle$ and first-excited $A_1|j = 3/2, \Omega' = 3/2\rangle$ states of the heteronuclear ion YX^+ ($g_{tot} = g_{j\Omega} + g_{j\Omega'} = 4$; ρ is the impact parameter; $\chi_{j\Omega,q\rho}(R)$ and $\chi_{j\Omega',q'\rho}(R)$ are the radial wave functions of the nuclei (normalized to a momentum δ function):

$$\chi_{j\Omega,q\rho}(R) = \left(\frac{2}{\pi}\right)^{1/2} \left(\frac{q}{q_{j\Omega}(R)}\right)^{1/2} \frac{1}{R} \cos\left[\int_b^R q_{j\Omega}(R)\,dR - \frac{\pi}{4}\right],$$

$$q_{j\Omega}(R) = \left[2\mu\left(E + \varepsilon_{nl} - U_{j\Omega,nl}(R) - E\frac{\rho^2}{R^2}\right)\right]^{1/2}, \quad (5)$$

which describe their relative motion with kinetic energies $E = q^2/2\mu$ and $E' = q'^2/2\mu = E + \varepsilon_{nl} - \varepsilon_{n'l'}$ as $R \to \infty$ ($q = q_{j\Omega}(\infty) = \mu v_E$ is the momentum v_E the velocity, b the left-hand turning point) in the Rydberg adiabatic[1)] terms $U_{j\Omega,nl}(R) = U_{j\Omega}(R) + \varepsilon_{nl}$ and $U_{j\Omega',n'l'}(R) = U_{j\Omega'}(R) + \varepsilon_{n'l'}$ of the $YX^+ + e^-$ quasimolecule (see Fig. 1). Note that the energy origin is chosen in (5) such that $U_{j\Omega}(\infty) = U_{j\Omega'}(\infty) = 0$.

In the calculation of the perturbation matrix elements, the electron wave functions $\varphi_{j\Omega}(\mathbf{r}_\varkappa, \mathbf{R})$ and $\varphi_{j\Omega'}(\mathbf{r}_\varkappa, \mathbf{R})$ of heavy heteronuclear ions of noble gases YX^+ (e.g, $HeAr^+$, $HeKr^+$, $HeXe^+$, $ArKr^+$, $ArXe^+$ and others, for which a Hund bond of type c is realized, see Refs. 6 and 7) can be expressed in the zeroth approximation, even at small internuclear distances, in terms of corresponding wave functions $\varphi_\varkappa(\mathbf{r}_\varkappa)$ and $\varphi_{j\Omega'}(\mathbf{r}_\varkappa)$ of the electron $n_0 p^5(^2P_{3/2})$ shell of the atomic ion X^+. This is permissible, since the electrostatic interaction of the colliding particles $X^+(^2P_{3/2})$ and $Y(^1S_0)$ is determined by the splitting $\Delta U^{(j)}_{\Omega\Omega'}(R) = U_{j\Omega'}(R) - U_{j\Omega}(R)$ of the terms $X|j = 3/2, \Omega = 1/2\rangle$ and $A_1|j = 3/2, \Omega' = 3/2\rangle$, which is much smaller (see Fig. 1) than the value $\Delta_{jj'}$ of the spin-orbit splitting of the $^2P_{3/2}$ and $^2P_{1/2}$ states of the ground electron shell of the X^+ ion [for $HeXe^+$, for example, we have $\Delta_{jj'} = 1.3$ eV and $0 \leqslant \Delta U^{(j)}_{\Omega\Omega'} \lesssim 0.1$ eV (Refs. 6 and 7)].

We recognize also that the perturbation operator $V = \Sigma^5_{\varkappa=1} |\mathbf{r}_e - \mathbf{r}_\varkappa|^{-1}$ does not depend on the internuclear distance R, since the investigated process is due to interaction of the Rydberg electron e^- with only the internal $n_0 p^5(^2P_{3/2})$ electrons of the X^+ ion (the coordinates $\mathbf{r}_\varkappa$ are measured from the X^+ nucleus), and its interaction with the electrons of the filled $^1S_0$ shell of the Y atom can be disregarded. This allows us to calculate the nuclear matrix element independently of the electronic ones. As a result we arrive, with the aid of the Landau-Zener method,[5] at the expression

$$|\langle \chi_{j\Omega,q\rho}(R) | \chi_{j\Omega',q'\rho}(R) \rangle|^2 = \frac{2q' \cos^2\left\{\int_b^{R_\omega} q_{j\Omega}(R)\,dR - \int_{b'}^{R_\omega} q_{j\Omega'}(R)\,dR + \pi/4\right\}}{\pi\mu\Delta F(R_\omega)\left[1 - U_{j\Omega}(R_\omega)/E - \rho^2/R_\omega^2\right]^{1/2}}, \quad (6)$$

where $\Delta F(R_\omega) = |dU_{j\Omega}/dR - dU_{j\Omega'}/dR|_{R^\omega}$ is the difference between the slopes of the considered terms of the $YX^+ + e^-$ quasimolecule at their intersection point R_ω determined from the condition $U_{j\Omega,nl}(R_\omega) = U_{j\Omega,n'l'}(R_\omega)$ or $\omega = \varepsilon_{nl} - \varepsilon_{n'l'} = \Delta U^{(j)}_{\Omega\Omega'}(R_\omega)$, see Fig. 1). It follows from the available quantum-mechanical calculations[6,7] that the power-law approximation $\Delta U^{(j)}_{\Omega\Omega'}(R) = A/R^\nu$ can be used for the energy splitting of the lower electronic terms $U_{j\Omega}(R)$ and $U_{j\Omega'}(R)$ of the heteronuclear YX^+ noble-gas ions (which have different components $\Omega = 1/2$ and $\Omega' = 3/2$ of the total electron angular momentum $j = 3/2$ along the internuclear axis $\mathbf{R}$) in a wide range of distances R. The expression for the radius R_ω of the crossing of the Rydberg terms takes then the following simple form:

$$R_\omega = (A/\omega)^{1/\nu}, \quad (7)$$

where $\omega = |\Delta n_{\text{eff}}|/n^3_{\text{eff}}$ is the frequency of the transition $nl \to n'l'$, and $\Delta n_{\text{eff}} = n' - \delta_{l'} - n + \delta_l$.

2.2 Electronic matrix elements

In the calculation of the electronic matrix elements $\langle j\Omega | \langle nlm | V | n'l'm' \rangle \times | j\Omega' \rangle$ we change with the aid of Clebsch-Gordan coefficients (see Refs. 5 and 8) from the wave functions $\varphi_{j\Omega}(\mathbf{r}_\varkappa)$ of the electron $n_0 p^5(^2P_{3/2})$ shell of the X^+ ion, expressed in the j-j coupling approximation and corresponding to the $j\Omega$ representation ($j = 3/2$, $\Omega = 1/2$ and $\Omega' = 3/2$), to the coordinate $\varphi_{n_0 L\Lambda}(\mathbf{r}_\varkappa)$ and spin $\eta_{S\mu}$ wave functions of the $LS\Lambda\mu$ representation ($L = 1$, $\Lambda = 0.1$ and $S = 1/2$; $\mu = \pm 1/2$ are the orbital momentum, the spin, and their possible projections on the intranuclear axis R for the $n_0 p^5$ shell)

$$\varphi_{j,\Omega=1/2} = \left(\frac{2}{3}\right)^{1/2} \varphi_{n_0 L,\Lambda=0}(\mathbf{r}_\varkappa)\,\eta_{S,\mu=1/2} + \frac{1}{3^{1/2}} \varphi_{n_0 L,\Lambda'=1}(\mathbf{r}_\varkappa)\,\eta_{S,\mu=-1/2},$$

$$\varphi_{j,\Omega'=3/2} = \varphi_{n_0 L,\Lambda'=1}(\mathbf{r}_\varkappa)\,\eta_{S,\mu=1/2}.$$

As a result, recognizing that the perturbation operator V contains no spin variables, we get

$$\langle \varphi_{j\Omega} | \langle nlm | V | n'l'm' \rangle | \varphi_{j\Omega'} \rangle = (^2/_3)^{1/2} \langle \varphi_{n_0 L\Lambda} | \langle nlm | V | n'l'm' \rangle | \varphi_{n_0 L\Lambda'} \rangle. \quad (8)$$

In addition, we replace the interaction

$V=\Sigma_{\varkappa=1}^{5}|\mathbf{r}_e-\mathbf{r}_\varkappa|^{-1}$ of the Rydberg electron e^- with all the electrons $\mathbf{r}_\varkappa$ ($\varkappa=1$–5) of the n_0p^5 shell by the equivalent interaction $V_{ev}=-r_{ev}^{-1}=-|\mathbf{r}_e-\mathbf{r}_v|^{-1}$ of this electron e^- with a vacancy ($\mathbf{r}_v$) in the noble-gas atom $X(n_0p^5)$. Indeed, supplementing the inner n_0p^5 shell by one more equivalent electron ($\mathbf{r}_\varkappa,\varkappa=6$) with wave function $\varphi_{n_0\mathscr{L}_6\lambda_6}(\mathbf{r}_\varkappa)$ (where $\mathscr{L}_6=1$ and $\lambda_6=0.1$ are the orbital momentum of this electron and its projection on the z axis) and by a positively charged vacancy with the same wave function $\varphi_{n_0L\Lambda}(\mathbf{r}_v)$, we rewrite the interaction V in the form $V=V_0+V_{ev}$, where $V_0=\Sigma_{\varkappa=1}^{6}|\mathbf{r}_e-\mathbf{r}_\varkappa|^{-1}$. It is easily seen that the interaction V_0 does not contribute to the considered transition $|n_0, L=1, \Lambda=0\rangle\to|n_0, L=1, \Lambda'=1\rangle$, since the vacancy wave functions are orthogonal, $\langle\varphi_{n_0L\Lambda}(\mathbf{r}_v)|\varphi_{n_0L\Lambda'}(\mathbf{r}_v)\rangle=0$. Thus, the investigated transitions of the Rydberg electron are determined by its interaction $V_{ev}=-|\mathbf{r}_e-\mathbf{r}_v|^{-1}$ with the vacancy, which depends on the coordinates of the n_0p^6 shell. Such transitions $|L=1, \Lambda=0; nl\rangle\to|L=1,\Lambda'=1; n'l'\rangle$ in the YX^++e^- system can therefore be regarded as an $nl\to n'l'$ transition of a Rydberg electron and a $|\varphi_{n_0L\Lambda}(\mathbf{r}_v)\rangle\to|\varphi_{n_0L\Lambda'}(\mathbf{r}_v)\rangle$ transition of a one-electron vacancy in the YX^+ ion, disregarding the electrons from the closed n_0p^6 shell.

We show next that for the mechanism considered in this article, a collisional quenching due to $nl\to n'l'$ transition of a Rydberg electron with change $\Delta\Omega=\Omega'-\Omega=1$ of the projection of the total angular momentum $j=3/2$ of the inner electrons of the quasimolecular ion YX^+ (and with a corresponding change $\Delta\Lambda=1$ of the projection of the orbital momentum $L=1$), contributions to the perturbation matrix elements are made only by transitions with $\Delta l=0$, ± 2 and $\Delta m=-1$. In fact, using the known expansion of the operator $V_{ev}=-r_{ev}^{-1}=-|\mathbf{r}_e-\mathbf{r}_v|^{-1}$ in terms of the multipole moments $r_<^k/r_>^{k+1}$ (see, e.g., Ref. 8), and representing the wave functions of the outer electron

$$\psi_{nlm}(\mathbf{r}_e)=R_{nl}(r_e)\,Y_{lm}(\theta_{\mathbf{r}_e},\varphi_{\mathbf{r}_e})$$

and of the vacancy

$$\varphi_{n_0L\Lambda}(\mathbf{r}_v)=R_{n_0L}(r_v)\,Y_{L\Lambda}(\theta_{\mathbf{r}_v},\varphi_{\mathbf{r}_v})$$

(where $L=1$, $\Lambda=0$, $\Lambda'=1$) in the form of products of radial and angular parts, we get

$$\langle\varphi_{nL\Lambda}|\langle nlm|r_{ev}^{-1}|n'l'm'\rangle|\varphi_{n_0L\Lambda'}\rangle=\sum_{k=0}^{\infty}\sum_{q=-k}^{k}3(-1)^{m-k}\,i^{l'-l}$$
$$\times\begin{pmatrix}1&k&1\\0&0&0\end{pmatrix}\begin{pmatrix}1&k&1\\0&q&1\end{pmatrix}\begin{pmatrix}l&k&l'\\0&0&0\end{pmatrix}\begin{pmatrix}l&k&l'\\-m&-q&m'\end{pmatrix}$$
$$\times\langle R_{n_0L}(r_v)|\langle R_{nl}(r_e)|r_<^k/r_>^{k+1}|R_{n'l'}(r_e)\rangle|R_{n_0L}(r_v)\rangle.\quad(9)$$

Here $r_<=\min(r_e,r_v)$ and $r_>=\max(r_e,r_v)$. Since the $3j$-symbols in (9) differ from zero only if $k=2$ and $q=-1$, we obtain with the aid of (8) and (9) for the squared modulus of the electronic matrix element

$$\sum_{m,m'}|\langle\varphi_{j\Omega}|\langle nlm|r_{ev}^{-1}|n'l'm'\rangle|\varphi_{j\Omega'}\rangle|^2=\frac{2}{25}C_{ll'}I_{nl,n'l'}^2,\quad(10a)$$

$$I_{nl,n'l'}=\int_0^\infty r_v^2\,dr_v\,R_{n_0L}^2(r_v)\left\{\int_0^{r_v}\frac{r_e^2}{r_v^3}R_{nl}(r_e)R_{n'l'}(r_e)r_e^2\,dr_e\right.$$
$$\left.+\int_{r_v}^\infty\frac{r_v^2}{r_e^3}R_{nl}(r_e)R_{n'l'}(r_e)r_e^2\,dr_e\right\},\quad(10b)$$

$$C_{ll'}=(2l+1)(2l'+1)\begin{pmatrix}l&2&l'\\0&0&0\end{pmatrix}^2\sum_{m,m'}\begin{pmatrix}l&2&l'\\-m&-1&m'\end{pmatrix}^2.\quad(10c)$$

Here $I_{nl,n'l'}=I_{nl,n'l'}^{(1)}+I_{nl,n'l'}^{(2)}$ is a radial matrix element of the sum $r_<^2/r_>^3$ of the short-range $I_{nl,n'l'}^{(1)}$ and quadrupole (long-range) $I_{nl,n'l'}^{(2)}$ parts of the interaction of the electron e^- with the vacancy r_v. It follows from (10a) that the cross sections of the considered $nl\to n'l'$ transitions differ from zero only for $l'=l$ ($l\neq0$) and $l'=l\pm2$.

In the calculation of the short-range part of the radial matrix element, the Coulomb wave functions of the Rydberg electron e^- can be expanded in a series near the origin $r_e=0$ (see Ref. 9)

$$R_{nl}(r_e)=\frac{2^{l+1}r_e^l}{n_{eff}^{3/2}(2l+1)!}\left[1-\frac{2r_e}{2l+2}+\ldots\right],$$

since contributions to $I_{nl,n'l'}^{(1)}$ are made by distances $0\leqslant r_e\lesssim\langle r_v\rangle_{n_0p}\sim1$ a.u. that are small compared with the orbit radius $\langle r_e\rangle_{nl}\sim2n_{\rm eff}^2$ of the highly excited atom $X(nl)$. Substituting the first terms of this expansion for the wave functions of the initial $R_{nl}(r_e)$ and final $R_{n'l'}(r_e)$ states in Eq. (10b) for the matrix element $I_{nl,n'l'}^{(1)}$, and calculating the integral over dr_e, we arrive at the following result:

$$I_{nl,n'l'}^{(1)}=\frac{2^{l+l'+2}\langle r_v^{l+l'+2}\rangle_{n_0p}}{(n_{eff}n'_{eff})^{3/2}(2l+1)!(2l'+1)!(l+l'+5)},$$

$$\langle r_v^s\rangle_{n_0p}=\int_0^\infty r_v^sR_{n_0p}^2(r_v)r_v^2\,dr_v,\quad(11)$$

where $\langle r_v^s\rangle_{n_0p}$ is a vacancy multipole of order s in the ground shell of the X^+ ion with principal quantum number n_0 and with orbital momentum $L=1$.

The quadrupole part $I_{nl,n'l'}^{(2)}$ of a radial matrix element can be calculated using the quasiclassical approximation for the wave function of the outer electron

$$R_{nl}(r_e)=\left(\frac{2}{\pi n_{eff}^3}\right)^{1/2}\frac{\cos\Phi(r_e)}{r_ep^{1/2}(r_e)}$$

$$\Phi(r_e)=\int_{a_1}^{r_e}p(r_e)dr_e-\frac{\pi}{4},\qquad p(r_e)=\left(-2|\varepsilon_{nl}|+\frac{2}{r_e}-\frac{\tilde l^2}{r_e^2}\right)^{1/2},$$
$$(12)$$

where $\tilde l=l+1/2$, $a_{1,2}=n_{\rm eff}^2\pm n_{\rm eff}(n_{\rm eff}^2-\tilde l^2)^{1/2}$ are the left- and right-hand turning points. Within the scope of the Fourier-component method, expression (10b) for $I_{nl,n'l'}^{(2)}$ can be represented with the aid of (12) in the form

$$I_{nl,n'l'}^{(2)}=\frac{1}{\pi(n_{eff}n'_{eff})^{3/2}}\int_0^\infty r_v^4\,dr_v\,R_{n_0L}^2(r_v)\int_{\max\{r_v,a_1\}}^{a_2}\frac{\cos\Delta\Phi(r_e)}{r_e^3p(r_e)}dr_e,$$

$$\Delta\Phi=\omega\int_{a_1}^{r_e}\frac{dr_e}{p(r_e)}-\tilde l\Delta l\int_{a_1}^{r_e}\frac{dr_e}{r_e^2p(r_e)}.\quad(13)$$

For the most interesting case of $nl\to n'l$ transitions without change of the orbital momentum ($\Delta l=0$) we can obtain a simple analytic estimate of the quadrupole part of a radial matrix element by neglecting in (13) the weak dependence on the frequency $\omega=\Delta n_{\rm eff}/n_{\rm eff}^3$, i.e., assume that

$\cos\Delta\Phi \approx 1$, and set the lower limit of the integration over dr_e approximately equal to $r_v \approx a_1$, (where $a_1 \approx \tilde{l}^2/2$ for $l \ll n$). Changing over in (13) to the dimensionless variable $x_e = r_e/2n_{\rm eff}^2$ and recognizing that in this case the momentum of the Rydberg electron see (12)] is given by

$$p(x_e) = X^{1/2}(x_e)/n_{eff}x_e, \text{ where } X(x_e) = x_e - x_e^2 - \tilde{l}^2/4n_{eff}^2,$$

we obtain

$$I_{nl,\,n'l}^{(2)} \approx \frac{1}{4\pi n_{eff}^6}\int_0^\infty r_v{}^4 R_{n_0p}^2(r_v)\,dr_v \times \int_{a_1/2n_{eff}^2}^{a_2/2n_{eff}^2} \frac{dx_e}{x_e{}^2 X^{1/2}(x_e)} = \frac{\langle r_v{}^2\rangle_{n_0p}}{n_{eff}^3 \tilde{l}^3}. \qquad (14)$$

This equation yields a reasonable result, since a vacancy in the n_0p^5 shell of the considered noble-gas X^+ ion is located mainly in the vicinity of the left-hand turning point for a Rydberg electron. Thus, for example, for the most effective $np \to n'p$ transitions we have $a_1 = 1.125$ a.u., and the characteristic vacancy sizes are $\langle r_v\rangle_{n_0p} \approx 1.7$ a.u. for Ar^+ and $\langle r_v\rangle_{n_0p} \approx 2.3$ a.u. for Xe^+, i.e., $\langle r_v\rangle_{n_0p} \sim a_1 \ll \langle r_e\rangle_{nl}$ (where $\langle r_e\rangle_{nl} \sim 2n_{\rm eff}^2$ is the average radius of the Rydberg-electron orbit).

In the general case, for allowed transition with both $\Delta l = 0$ and $\Delta l = \pm 2$, it is convenient to write the exact result for the matrix element $I_{nl,n'l'}^{(2)}$, which takes into account the change of the frequency $\omega = \Delta n_{\rm eff}/n_{\rm eff}^3$, in the form

$$I_{nl,n'l'}^{(2)} = \frac{\langle r_v{}^2\rangle_{n_0p}}{n_{eff}^3\tilde{l}^3} f(\Delta n_{eff}, \Delta l);$$

$$f(\Delta n_{eff}, \Delta l) = \frac{1}{\langle r_v{}^2\rangle_{n_0p}}\int_0^\infty r_v{}^4 R_{n_0p}^2(r_v) F(r_v)\,dr_v, \qquad (15a)$$

$$F(r_v) = \frac{\tilde{l}^3}{4\pi n_{eff}^3}\int_{x_0}^{a_2/2n_{eff}^2} \frac{dx_e}{x_e{}^2 X^{1/2}(x_e)} \cos\Delta\Phi(x_e), \qquad (15b)$$

where $x_0 = \max\{a_1/2n_{\rm eff}^2, r_v/2n_{\rm eff}^2\}$. Calculation of the quasiclassical phase difference $\Delta\Phi$ defined by Eq. (13) leads to the expression

$$\Delta\Phi(x_e) = -\Delta n_{eff}\left\{\arcsin\left[\frac{1}{e_l}(1-2x_e)\right] + 2X^{1/2}(x_e) - \frac{\pi}{2}\right\} - \Delta l\left\{\arcsin\left[\frac{1}{e_l}\left(1 - \frac{\tilde{l}^2}{2n_{eff}^2 x_e}\right)\right] - \frac{\pi}{2}\right\}, \qquad (15c)$$

where $e_l = (1 - \tilde{l}^2/n_{\rm eff}^2)^{1/2}$ is the eccentricity of the Rydberg-electron orbit.

It follows directly from this expression that for $nl \to n'l$ transitions with $\Delta l = 0$ the value of $\cos\Delta\Phi(x_e)$ is practically independent of $x_e = r_e/2n_{\rm eff}^2$ in the distance region $x_e \sim \tilde{l}^2/n_{\rm eff}^2$ that makes the main contribution to the integral in (15b). For $nl \to n'l$ quadrupole transitions with change of only the principal quantum number, the value of $f(\Delta n_{\rm eff}, \Delta l = 0$ therefore depends weakly on $\Delta n_{\rm eff}$ in the entire range of transition frequencies $\omega = \Delta n_{\rm eff}/n_{\rm eff}^3$ considered in this paper, and can be determined with sufficient accuracy (as follows from the results of a numerical calculation for different values of $\Delta n_{\rm eff}$) for $\Delta n_{\rm eff} = 0$, when $\cos\Delta\Phi = 1$. This leads to the following analytic result for the function (15b):

$$F(r_v) = \frac{1}{2} - \frac{1}{\pi}\arcsin\left[\frac{1}{e_l}\left(1 - \frac{\tilde{l}^2}{r_v}\right)\right] - \frac{2X^{1/2}(r_v/2n_{eff}^2)}{\pi n_{eff}^2 \tilde{l}^2 r_v}. \qquad (16)$$

Thus, calculation of the quadrupole matrix element $I_{nl,n'l}^{(2)}$ of the transition reduces to averaging of $r_v^2 F(r_v)$ [see Eqs. (15a) and (16)] over the Hartree-Fock wave functions $R_{n_0p}^2(r_v)$ of the electron n_0p^5 shell of the ion X^+. The result $f(\Delta n_{\rm eff}, \Delta l = 0)$ of this averaging is found to be weakly dependent on the principal quantum number n, i.e., reduces in practice to a certain constant coefficient $\langle f(\Delta n_{\rm eff}, \Delta l = 0)\rangle_n = f_0$ ($f_0 = 0.46$ for Xe^+) in the equation for the radial matrix element of the $nl \to n'l$ transition [see (15a)]; this enables us to refine the simple estimate (14) obtained above.

For $nl \to n'l'$ transitions with orbital-momentum change $\Delta l = \pm 2$ the quadrupole matrix elements $I_{nl,n'l'}^{(2)}$ turn out to be significantly smaller [as follows from a numerical calculation of $f(\Delta n_{\rm eff}, \Delta l = \pm 2)$] than for transitions with $\Delta l = 0$. An estimate based on the approximation $\langle r_v\rangle_{n_0p} \sim a_1$ used above leads to the same conclusion, inasmuch as for $\omega = 0$ the integral (15a) vanishes identically for transitions with $\Delta l = \pm 2$, owing to the oscillations of $\cos\Delta\Phi(x_e)$ [see (15c)]. In the following actual calculations of the cross sections of $nl \to n'l'$ transitions with $\Delta l = \pm 2$ and $l = 0, 1, 2$ we shall therefore neglect the contribution of the quadrupole (long-range, $r_v < r_e < a_2$) part of the interaction V compared with the contribution [see (11)] of the short-range ($0 \leqslant r_e \lesssim r_v$) part of the interaction.

Thus, with the aid of Eqs. (8), (10), (11), and (15) the squared modulus of the electronic transition matrix element

$$|X, j={}^3/_2, \Omega={}^1/_2;\ nl\rangle \to |A_1, j={}^3/_2, \Omega'={}^3/_2;\ n'l'\rangle$$

can be represented in the form

$$\sum_{m,m'} |\langle\varphi_{j\Omega}|\langle nlm|r_{ev}{}^{-1}|n'l'm'\rangle|\varphi_{j\Omega'}\rangle|^2 = \frac{2}{25}\frac{C_{ll'}Q_{ll'}^2}{n_{eff}^3 n_{eff}'^3}, \qquad (17a)$$

$$Q_{ll'} \approx \frac{2^{l+l'+2}\langle r_v{}^{l+l'+2}\rangle_{n_0p}}{(2l+1)!(2l'+1)!(l+l'+5)} + f_0\frac{\langle r_v{}^2\rangle_{n_0p}}{(l+{}^1/_2)^3}\delta_{ll'}. \qquad (17b)$$

2.3. Cross sections and rate constants of Rydberg-level quenching

We obtain an analytic result for the cross section of the investigated process (2b) using Eqs. (17) and (6) respectively for the electronic and nuclear matrix elements, after integrating the initial expression (4) over the impact parameters $2\pi\rho d\rho$ of the colliding particles $X(nl)$ and Y. We replace the quantity $\cos^2\{S_0(\rho) + \pi/4\}$ in (6), which oscillates rapidly when ρ is changed, by its mean value 1/2. As a result we have

$$\sigma_{nl,n'l'}^{\Omega\Omega'} = \frac{8\pi^2\gamma_{ll'}R_\omega{}^2[1-U_{j\Omega}(R_\omega)/E]^{1/2}}{25(2l+1)n_{eff}^3 n_{eff}'^3 v_E \Delta F(R_\omega)} = \frac{8\pi^2 A^{3/\nu}\gamma_{ll'}[1-U_{j\Omega}(R_\omega)/E]^{1/2}}{25(2l+1)\nu v_E n_{eff}^{3(1-3/\nu)}|\Delta n_{eff}|^{1+3/\nu}}, \qquad (18)$$

where $\gamma_{ll'} = C_{ll'}Q_{ll'}^2$ is a parameter that characterizes the magnitude of the interaction. The second equation in (18) pertains to the case of a power-law approximation of the term-splitting energy of the heteronuclear ion YX^+ (see Sec. 2.1), when the intersection radius R_ω of the corresponding Rydberg terms of the $YX^+ + e^-$ quasimolecule is determined by the simple exprsssion (7). An actual analysis shows that the most effective are the $np \to n'p$ transitions with $l' = l = 1$. Thus, for example, for $Xe(^2P_{3/2}, nl) + He(^1S_0)$ collisions we obtain with the aid of Eqs. (10c) and (17b) $\gamma_{pp} = 2.04$ a.u., and the contribution from all other transitions with $l' \neq 1$ and $l \neq 1$ decreases rapidly with increase of the orbital momentum l or l' and turns out to be smaller by more than an order of magnitude than γ_{pp} ($\gamma_{sd} = \gamma_{ds} = 0.083$ a.u., $\gamma_{dd} = 0.019$ a.u., etc.). The multipole moments $\langle r_v^2 \rangle_{5p} = 5.6$ a.u., $\langle r_v^4 \rangle_{5p} = 33.9$ a.u., and others, which determine the radial integrals $Q_{ll'}$, [see (11) and (17b)], are calculated in this case by using the Hartree-Fock wave functions $R_{5p}(r_v)$ of the electron $5p^5$ shell of the Xe^+ ion (see Ref. 10).

For the total rate $W_n^Y = K_n^Y N_Y$ of the collisional quenching of a specified energy level n of a highly excited atom $X(n)$ by atoms of a buffer gas Y we obtain after integrating (18) over the Maxwellian distribution of the velocities $v_E = (2E/\mu)^{1/2}$ of the nuclei, summing over all the final $n'l'$ levels with $n' < n$, and averaging over the initial l-sublevels

$$W_n{}^Y = \sum_{n'<n} \sum_{ll'} \frac{2l+1}{n^2} \langle v_E \sigma_{nl,n'l'}^{\Omega,\Omega'} \rangle N_Y = \frac{8\pi^2 \gamma A^{3/\nu} N_Y}{25\nu n^{5-9/\nu}} \sum_{\Delta n<0} \frac{a[U_{j\Omega}(R_\omega)/T]}{|\Delta n|^{1+3/\nu}} \qquad (19)$$

$$a = \frac{2}{\pi^{1/2}} \int_{\max\{0,U(R_\omega)/T\}}^{\infty} d\left(\frac{E}{T}\right) \exp\left(-\frac{E}{T}\right) \left[\frac{E}{T} - \frac{U_{j\Omega}(R_\omega)}{T}\right]^{1/2} = \frac{2}{\pi^{1/2}} \exp\left[-\frac{U_{j\Omega}(R_\omega)}{T}\right] \Gamma\left(\frac{3}{2}, \frac{|U_{j\Omega}(R_\omega)|}{T}\right).$$

Account is taken in this equation of all the effective $np \to n'p$ transitions in the $X(nl)$ atom, with frequencies $\omega = \Delta n / n^3$ (where $\Delta n = n' - n = -1, -2, -3, \ldots$), which lead to quenching of the Rydberg electron, so that $\gamma \approx \gamma_{pp} = C_{11} Q_{11}^2$ [see (10c) and (17b)], and the following notation is introduced: N_Y is the density of the Y buffer-gas atoms, T is the gas temperature, $a[U_{j\Omega}(R_\omega)/T]$ is a factor that takes into account the change of the relative velocity of the colliding nuclei X^+ and Y as they move in the potential well of the principal term $X\,|j = 3/2, \Omega = 1/2\rangle$ expressed in terms of the incomplete gamma function $\Gamma(3/2, z)$ (see Ref. 11). For the heteronuclear YX^+ noble-gas ions investigated here the well depth for this term is ~ 0.03–0.1 eV (see Refs. 6 and 7), i.e., of the order of the thermal energy $E \sim T \sim 0.03$–0.05 eV of the atoms $X(n)$ and Y. Therefore the value of $a[U_{j\Omega}(R_\omega)/T]$ changes little over the entire investigated range of the frequencies ω and can be replaced by an appropriate mean value. Then, summing over Δn in (19), we get ultimately

$$W_n{}^Y = K_n{}^Y N_Y = \frac{8\pi^2 \gamma A^{3/\nu} \zeta(1+3/\nu)}{25\nu n^{5-9/\nu}} a\left[\overline{\frac{U_{j\Omega}(R_\omega)}{T}}\right] N_Y, \qquad (20)$$

where $\zeta(z) = \Sigma_{k=1}^{\infty} k^{-z}$ is the Riemann zeta function (see Ref. 11). For the specific case of $Xe(n) + He$ collisions

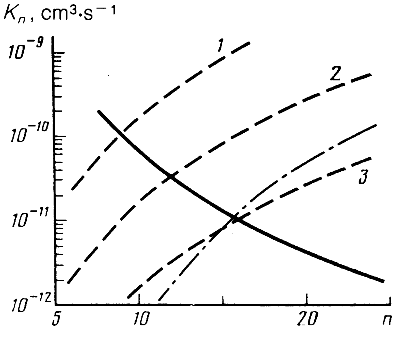

FIG. 2. Comparison of the probabilities of the inelastic quenching of $Xe(n)$ Rydberg levels by He buffer-gas atoms and by free electrons of an He/Xe mixture plasma ($N_{Xe} < N_{He}$). Solid curve—quenching rate constants $K_n = \Sigma_{n',l',l}[(2l+1)/n^2] K_{nl,n'l'} = \langle v_E \sigma_n^{He} \rangle$ for the mechanism considered in the paper [calculated using Eq. (20)]. Dash-dot curve—rate constants K_n for quenching by the competing Fermi mechanism [calculated using Eq. (24) of Ref. 12]. Dashed curves—relative probabilities $W_n^e / N_{He} = \langle v_e \sigma_n^e \rangle \alpha$ of quenching highly excited levels by electron impact (calculated using Eq. (10.17) of Ref. 14) with $T_e = 0.2$ eV) for different degrees of plasma ionization $\alpha = N_e / N_{He} = 10^{-6}$, 10^{-7} and 10^{-8} (curves 1, 2, and 3, respectively).

(when $A = 1.6 \cdot 10^3$ a.u.; $\nu = 7.8$; $\gamma = 2.04$ a.u.; $a[\overline{U_{j\Omega}(R_\omega)}/T] \approx 1.5$ for $T = 300$–600 K) we obtain from (20) the simple estimate K_n^{He} [cm³·s⁻¹] $\approx 3.10^{-7}\, n^{-3.85}$ for the total n-level energy quenching rate constant.

As seen from Fig. 2, quenching of $Xe(n)$ Rydberg atoms in inelastic collisions with He atoms via the mechanism proposed here is substantially more effective than via the traditional Fermi mechanism in the region $n \lesssim 15$–20 (the quenching rates for this case were calculated with the aid of the analytic equations (20) and (24) of Ref. 12). The probabilities $W_n^e = \langle v_e \sigma_n^e \rangle N_e$ of quenching highly excited levels of $Xe(n)$ atoms by electron impact (see Refs. 13 and 14) also turn out to be lower than the probabilities $W_n^{He} = \langle V_E \sigma_n^{He} \rangle N_{He}$ calculated here [see (19)] of quenching the levels in collisions with He buffer-gas atoms for quantum numbers $n \lesssim 8$, 12, and 16 respectively at plasma ionization degrees $\alpha = 10^{-6}$, 10^{-7}, and 10^{-8} (see Fig. 2). It follows hence that the considered quenching Rydberg levels of $X(n)$ atoms, via scattering of the perturbing Y atom by the ionic core X^+, should influence strongly the rates of three-particle recombination of electrons with atomic X^+ ions in a low-temperature plasma of an X/Y noble-gas mixture.

3. KINETIC MODEL OF RECOMBINATION

Consider a three-particle recombination model in which the recombining electron relaxes over the highly excited levels of the $X(n)$ atom both by collisions with free electrons (1) and by the above mechanism of excitation quenching by collisions with buffer-gas atoms (2). Within the framework of the modified diffusion approximation,[3] just as in the case of recombination of atomic X^+ ions with electrons in the gas proper,[4] we replace the discrete energy-level spectrum by a quasi-continuum and represent the recombination flux Γ in the stationary drain regime in the form

$$\Gamma = N_0(\varepsilon) D_e(\varepsilon) dy(\varepsilon)/d\varepsilon - W_Y(\varepsilon) y(\varepsilon),$$
$$y(\varepsilon) = N(\varepsilon)/N_0(\varepsilon), \tag{21}$$
$$D_e(\varepsilon) = 4(2\pi)^{1/2}\Lambda_C \varepsilon N_e/3T_e^{1/2},$$
$$N_0(\varepsilon) = \pi^{3/2} e^{\varepsilon/T_e} N_e N_{X^+}/4T_e^{3/2}\varepsilon^{5/2}.$$

Here $N(\varepsilon) = N_n |dn(\varepsilon)/d\varepsilon|$ is the electron density in the quasicontinuum with binding energy $\varepsilon \equiv |\varepsilon_n| = 1/2n^2$; $N_0(\varepsilon)$ is their equilibrium density determined by the Saha-Boltzmann equation; $D_e(\varepsilon)$ is the diffusion coefficient of a weakly bound electron in collisions with free plasma electrons; Λ_C is the Coulomb logarithm (see Refs. 1 and 3); N_e and N_{X^+} are the densities of the electrons and atomic ions in the plasma; $W_Y(\varepsilon)$ is the probability of quenching a highly excited electron of the $X(n)$ atom by atoms of the buffer gas Y, expressed in the quasicontinuum and connected with W_n^Y [see (20)] by the simple relation

$$W_Y(\varepsilon) = W^Y_{n(\varepsilon)} \left|\frac{d\varepsilon}{dn}\right|$$
$$= \frac{2^{7-9/2\nu}\pi^2}{25\nu}\zeta\left(1+\frac{3}{\nu}\right)\gamma A^{3/\nu} a\left[\frac{\overline{U_{j\Omega}(R_\omega)}}{T}\right]\varepsilon^{(8\nu-9)/2\nu} N_Y.$$

The solution of Eq. (21) with allowance for the boundary condition $y(0) = 1$ [zero energy $\varepsilon = 0$ corresponds to the ionization limit of the atom $X(n)$] is

$$y(\varepsilon) = \Pi(\varepsilon)\left[1 - \Gamma\int_0^{\varepsilon} \frac{d\varepsilon'}{N_0(\varepsilon')D_e(\varepsilon')\Pi(\varepsilon')}\right],$$
$$\Pi(\varepsilon) = \exp\left[\int_0^{\varepsilon}\frac{d\varepsilon' W_Y(\varepsilon')}{D_e(\varepsilon')}\right] = \exp\left[0{,}2\left(\frac{\varepsilon}{\varepsilon_Y}\right)^{(8\nu-9)/2\nu}\right], \tag{22}$$

$$\varepsilon_Y(T_e,\alpha) = \frac{1}{2}\left[\frac{5(8\nu-9)\Lambda_C\alpha}{3(2\pi)^{3/2}\gamma A^{3/\nu} a(\overline{U_{j\Omega}(R_\omega)}/T)\zeta(1+3/\nu)T_e^{1/2}}\right]^{2\nu/(8\nu-9)}, \tag{23}$$

where $\alpha = N_e/N_Y$ is the degree of ionization of the noble-gas X/Y mixture plasma (with $N_X \ll N_Y$). The electron binding energy ε_Y divides the discrete spectrum of the $X(n)$ atom into two regions. In the first ($\varepsilon \equiv |\varepsilon_n| \lesssim \varepsilon_Y$) the recombination flux Γ is determined mainly by diffusion over the highly excited levels in collisions with electrons. In the second ($\varepsilon \equiv |\varepsilon_n| \gtrsim \varepsilon_Y$) it is determined by the mechanism considered here for quenching by the neutral Y particles. For $\varepsilon < \varepsilon_Y$, in fact, we have $\Pi(\varepsilon) \approx 1$ and the solution (22) for $y(\varepsilon)$ coincides with the Gurevich-Pitaevskiĭ result,[1] i.e., quenching by the buffer gas is negligible. On the contrary, for $\varepsilon > \varepsilon_Y$ we have from (22) $\Pi(\varepsilon) \gg 1$, i.e., collisions with free electrons can be neglected. The transition from the "diffusion" region $\varepsilon < \varepsilon_Y$ to the "predominant drain" region $\varepsilon > \varepsilon_Y$ occurs in a narrow vicinity of the point ε_Y [see (23)]. We can therefore put $\Pi(\varepsilon) = 1$ for all $\varepsilon < \varepsilon_Y$ and, imposing a second boundary condition $y(\varepsilon_Y) = 0$, determine from (22) the recombination flux Γ and the coefficient $\beta_{eY} = \Gamma/N_e^3$ of three-particle recombination for the considered kinetic model:

$$\beta_{eY}(T_e, \alpha) = \beta_e(T_e)\xi^{-1}(\varepsilon_Y/T_e),$$
$$\beta_e = \frac{4\pi(2\pi)^{1/2}\Lambda_C}{9T_e^{9/2}}, \quad \xi(x) = \frac{4}{3\pi^{1/2}}\int_0^x t^{3/2}e^{-t}\,dt. \tag{24}$$

Here β_e is the three-particle recombination coefficient and is governed only by electron-electron collision, while $\beta[\text{cm}^6\cdot\text{s}^{-1}] = 8.8\cdot 10^{-27}\cdot\Lambda_C\ (T_e\ [\text{eV}])^{-9/2}$ (see Refs. 1 and 3). The quantity $\xi^{-1}(\varepsilon_Y/T_e)$ has the meaning of the coefficient of recombination-rate growth ($0 \leqslant \xi \leqslant 1$) due to quenching of Rydberg levels of $X(n)$ atoms by buffer-gas (Y) atoms, and is expressed in terms of an incomplete gamma function $\xi(x) = (\frac{4}{3}\pi^{1/2})\gamma(5/2, x)$ (see Ref. 11).

It follows directly from (24) that at electron temperatures $T_e \sim \varepsilon_Y$ [see (23)] the recombination regime changes from a steep decrease with increase of T_e at low temperatures

$$\Gamma \propto N_e^3 \cdot T_e^{-9/2} \text{ for } T_e \ll \varepsilon_Y,\ \xi(\varepsilon_Y/T_e) \approx 1, \tag{25a}$$

to a substantially slower decrease with temperature and anomalously large $\beta_{eY} \gg \beta_e$ at high temperatures:

$$\Gamma \propto N_e^{(19\nu-27)/(8\nu-9)} N_Y^{5\nu/(8\nu-9)} T_e^{-[(27\nu-36)/2(8\nu-\nu)]},$$
$$T_e \gg \varepsilon_Y, \quad \xi(\varepsilon_Y/T_e) \approx \frac{8}{15\pi^{1/2}}\left(\frac{\varepsilon_Y}{T_e}\right)^{5/2}. \tag{25b}$$

For the specific case of electron recombination with atomic Xe^+ ions in an He/Xe plasma at high electron temperatures $T_e \gg \varepsilon_Y$ we obtain from (25b) $\Gamma \propto N_e^{2.3} \times N_{He}^{0.7} \times T_e^{-1.6}$. Owing to the effective quenching of the Rydberg levels of the $Xe(n)$ atoms in the buffer gas, the recombination coefficient $\beta_{e,He}$ increases then compared with the values of β_e obtained when only the usual "diffusion" recombination with free plasma electrons is taken into account. In particular, for a temperature $T_e \approx 0.2$ eV at ionization degrees $\alpha = N_e/N_{He} = 10^{-5}, 10^{-6}, 10^{-7}$ we obtain from (24) an increase of the recombination rate $\beta_{e,He}/\beta_e$ by 4, 10, 57, and 325 times, respectively. This points to the efficacy of the quenching mechanism proposed for highly excited atom levels by neutral particles in electron-ion recombination of mixed noble-gas plasma.

4. EXPERIMENTAL INVESTIGATION OF ELECTRON-ION RECOMBINATION IN AN He/Xe MIXTURE GAS-DISCHARGE PLASMA

The recombination coefficients of electrons with Xe^+ ions in He as a buffer were measured in the afterglow plasma of a pulsed discharge in an He/Xe mixture at pressures $P_{He} = 8$–136 Torr and $P_{Xe} = 0.01$–0.3 Torr, in a tube of radius $R_0 = 0.6$ cm. The electron density was determined from the plasma conductivity using the known relation for the discharge current

$$i(t) = 2\pi e E(t)\int_0^{R_0} b_e[E(t)/N_{He}(R,t)]N_e(R,t)R\,dR, \tag{26}$$

where b_e is the electron mobility and $E(t)$ is the experimental longitudinal-electric-field strength determined by the two-probe method. The electron density at the initial instant of the decay (i.e., at the instant $t = t_0$ of current-pulse termination) was varied in the range $N_e(R = 0, t_0) = 2.5 \times 10^{11}$–

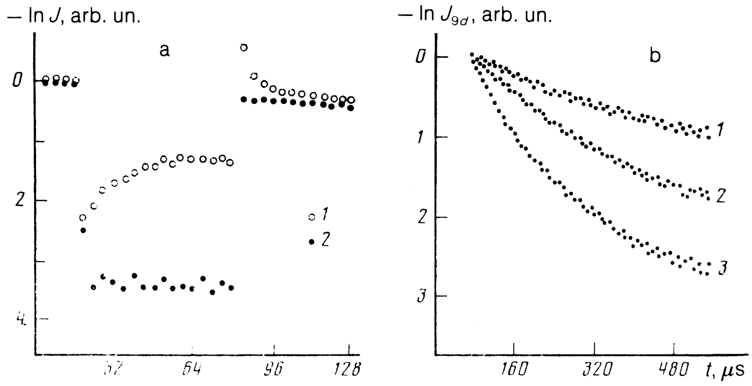

FIG. 3. Time dependences of the emission intensities of Xe-atom spectral lines in the afterglow of a xenon–helium plasma ($P_{He} = 100$ Torr, $P_{Xe} = 0.05$ Torr); a—pulsed heating of the electrons at $T_e(E/N_{He}) = 1000$ K and $N_e(t_0) = 2.5\times10^{11}$ cm^{-3} (curve 1 corresponds to the $7p\to6s$ ($\lambda = 467.1$ nm) line, curve 2 to the $9d\to6p$ ($\lambda = 593.4$ nm); b—stationary preheating of the electrons (curves 1, 2, and 3 correspond to a decrease, with time, of the $4d\to6p$ line intensity at electron temperatures $T_e = 540$ K (1), 640 K(2) and 1100 K (3)) for $T = 540$ K and $N_e(t_0) = 1.2\times10^{12}$ cm^{-3}.

1.5×10^{12} cm^{-3}, and the gas temperature was $T(R=0, t_0) = 350$–600 K, depending on the discharge power. The radial distributions and the values of $N_e(R,t)$, $N_{He}(R,t)$ and $T(R,t)$ were determined during the current pulse ($t<t_0$) by numerically solving a system comprising the gas heat-conduction equation, the equation for the densities of the electrons and metastable Xe[$6s(^3P_2)$] ions, and Eq. (1). The coefficients in these equations and the $b_e(E/N_{He})$ dependence were obtained from numerical calculations of the Boltzmann equation for the electron distribution function.

After the termination of the current pulse, the electrons cool off rapidly by elastic collisions with the He atoms, with characteristic times $\tau_e \lesssim 2\times10^{-6}$ s at $P_{He} \gtrsim 10$ Torr. The plasma decays therefore in the absence of a field at $T_e \approx T$. Estimates of the efficiencies of the internal electron-heating by recombination and by ionization in pair collisions of the metastable Xe[$6s(^3P_2)$] atoms have shown that these processes raise the electron temperature T_e by not more than 20 K.

The dependence of the recombination rate on the electron temperature in the range $T_e \approx 400$–2500 K were investigated by an electron "preheating" method based on the possibility, demonstrated in Ref. 15, of forming a specified longitudinal electric field $E(t)$ in a decaying plasma. The characteristic electron temperatures were determined in this case from the dependence, calculated for an He/Xe plasma, of the average electron energy on the parameter E/N_{He}.

The time dependences of the intensities $J_\lambda(t)$ of the spectral lines emitted by excitation of xenon atoms from the 9s-, (6–8) p, 6p′, (9–11) d, (4–7) f states were recorded in the experiments by the multichannel photon-counting method. The light beams were gathered from a narrow central region of the discharge tube. The measurements of $J_\lambda(t)$ were performed in a plasma in the absence of a field as well as for pulsed (see Fig. 3a for the $9d\to6p$ and $7p\to6s$ line intensities) and stationary (see Fig. 3b) preheating of the electrons in the afterglow stage.

The reported experimental results pertain to conditions when in the initial stage of the plasma decay one can neglect the population of all the excited levels of the Xe* atoms by recombination of the Xe_2^+ ions with electrons, compared with recombination of the atomic Xe^+ ions. The influence of dissociative recombination of the Xe_2^+ atoms could be monitored directly in the experiments because the differences between the indicated Xe* atom production are distinctly seen on the $J_\lambda(t)$ curves for pulsed preheating of the electrons (see Fig. 3a). Thus, for example, for the $7p\to6s$ line the influence of dissociative recombination was manifested by an increase of the intensity $J_\lambda(t)$ within the limits of the heating pulse. The reason is that the density $N_{Xe_2^+}(t)$ of the molecular ions increases on account of the decreased dissociative-recombination rate (as a result of the rise of T_e), while the rate of their formation by the conversion processes $Xe^+ + Xe + He \to Xe_2^+ + He$ remains unchanged. We note that this behavior was observed only during the final stage of the plasma decay (i.e., at low values of the electron density $N_e(t)$ for the 6p, 6p′, and 7p levels of the Xe* atoms with binding energy $|\varepsilon_{br}| > \mathscr{E}_0$, (where $\mathscr{E}_0$ is the binding energy of the molecular Xe_2^+ ion in the ground vibrational state $v=0$). These peculiarities of the behavior of $J_\lambda(t)$ do not occur for all the higher nl levels investigated here, with binding energy $|\varepsilon_{nl}| < \mathscr{E}_0 = 1$ eV [see, e.g., Fig. 3a for the radiation intensity $J_{9d}(t)$ of the 9d level ($\lambda = 593.4$ nm)]. This indicates incontrovertibly that the highly excited levels of the Xe(nl) atoms are populated by three-particle recombination of the electrons e^- with the atomic ions Xe^+.

From the foregoing data we determined the temperature dependences of the emission intensities $J_\lambda(T_e)$ of the spectral lines that start out from the recombination-populated levels of the Xe* atoms. As seen from Fig. 4a with the $J_{9d}(T_e)$ dependence of the emission of the $9d\to6p$ line as an example,[2)] the abrupt decrease $J_\lambda(T_e) \propto T_e^{-9/2}$ of the curves with temperature in the low-temperature region gives way to a much slower dependence at $T_e \gtrsim 800$ K. Measurements at various concentrations of the buffer gas He show that in this temperature region, $T_e \gtrsim 800$ K, the relative radiation intensities $J_\lambda(T_e, N_e, N_{He})/J_\lambda(T_e = 400\ \mathrm{K}, N_e)$ are approximately inversely proportional to the ionization degree $\alpha = N_e/N_{He}$.

For a more detailed experimental investigation of this effect, we have measured the absolute values of the recombination coefficients $\beta_{e,He}(T_e)$ of the electrons e^- with the atomic ions Xe^+. The values of $\beta_{e,He}$ were determined from data on the decrease of the intensities $J_{9d}(t)$ with time in a plasma in the absence of a field and with stationary heating of the electrons, using the following system of charged particle balance equations ($N_e = [Xe^+] + [Xe_2^+]$):

$$d[Xe^+]/dt = -\beta_{e,He}[Xe^+]N_e^2 - (B + \tau_D^{-1})[Xe^+], \qquad (27)$$

$$d[Xe_2^+]/dt = -(\alpha_{d,r}N_e + \tau_D^{-1})[Xe_2^+] + B[Xe^+]. \qquad (28)$$

Here $B = b$ [Xe] [He] is the reciprocal $Xe^+ + Xe + He \to Xe_2^+ + He$ conversion time in three-particle collisions ($b \approx 10^{-31}$ cm$^6\cdot$s^{-1}, Ref. 16), τ_D is the ambipolar dif-

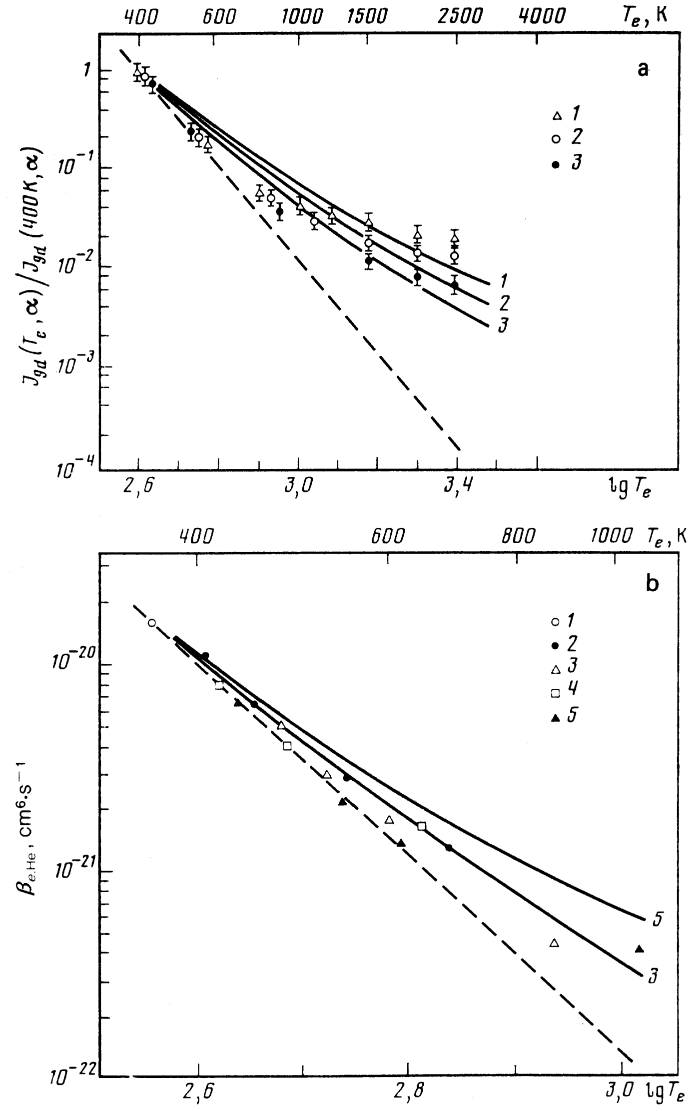

FIG. 4. Dependences of the relative intensities of the emission of the $9d \to 6p$ line of the Xe atom, and of the coefficient of three-particle recombination of electrons with Xe^+ ions, on the temperature T_e at different degrees of ionization α of a xenon–helium plasma: a—relative intensities $J_{9d}(T_e,\alpha)/J_{9d}(T_e = 400\ K)$ at pressures $P_{He} = 50$ Torr, $P_{Xe} = 0.05$ Torr and at the electron densities $N_e(t_0) = 10^{11}\ cm^{-3}$ (1), $2\times10^{11}\ cm^{-3}$ (2) and $4\times10^{11}\ cm^{-3}$ (3); b—recombination coefficients at pressures $P_{Xe} = 0.05$ Torr and $P_{He} = 55$ Torr (1, 2, 3), $P_{He} = 100$ Torr (4, 5) and the electron densities $N_e(t_0) = 2.5\times10^{11}$ (1), $6\times10^{11}\ cm^{-3}$ (2,4) and $1.2\times10^{12}\ cm^{-3}$ (3,5). Solid curves—corresponding calculations using Eqs. (23) and (24) of the present paper. Dashed curve—calculation using the equation $\beta_e(T_e)$ [cm^6·s^{-1}] $8.8\times10^{-27}\ \Lambda_C(T_e\ [eV])^{-9/2}$ (for $\Lambda_C = 0.25$) in which only collisions with electrons are taken into account.

fusion time, $\beta_{e,He}$ is the three-particle recombination coefficient of the atomic ions Xe^+, and $\alpha_{d,r}$ is the rate of dissociative recombination of the molecular ions Xe_2^+. No account is taken in these equations of ionization produced by two-particle collisions of metastable Xe(6s) atoms. Measurements by the absorption method show that the density of the Xe(6s) atoms did not exceed $10^{11}\ cm^{-3}$, so that they play no noticeable role in the charged-particle balance during the initial plasma-decay stage. Since the characteristic time $\tau_{d,r} = (\alpha_{d,r}N_e)^{-1}$ of the dissociative recombination ($\alpha_{d,r} = 10^{-6}\ cm^3\cdot s^{-1}$, Ref. 17) is smaller by two orders than all other characteristic times of the processes in (27) and (28), we have for the densities of the molecular ions $[Xe_2^+] = B[Xe^+]/\alpha_{d,r}N_e$ (i.e., $[Xe_2^+] < 5\times10^8\ cm^{-3}$ for $P_{He} = 50$–100 Torr, $P_{Xe} = 0.05$ Torr and $N_e = 2\cdot10^{11}$–$1.5\cdot10^{12}\ cm^{-3}$). This means that $[Xe^+] \approx N_e$, and the density $[Xe_2^+] \approx B/\alpha_{d,r}$ of the molecular ions is independent of the time during the initial plasma-decay stage. We obtain then from the set (27), (28) for the decrease of electron density with time

$$[N_e(t_0)/N_e(t)]^2 = \exp(-2t/\tau) + \beta_{e,He}\tau N_e^2(t_0)[\exp(2t/\tau) - 1], \qquad (29)$$

$$J_{9d}(t) \propto \beta_{e,He}N_e^3(t), \quad \tau^{-1} = \tau_D^{-1} + B. \qquad (30)$$

It is evident from (29) and (30) that the decrease of $N_e(t)$ and of the radiation intensity $J_{9d}(t)$ is determined by two parameters, $\beta_{e,He}$ and τ. The time τ was obtained from the decrease of the intensity $J_{9d}(t)$ of the line in the far afterglow, when $N_e(t) \propto \exp(-t/\tau)$. The recombination coefficient $\beta_{e,He}$ was determined from a least-squares comparison of the solution (29), (30) with the experimental $J_{9d}(t)$ dependence of the intensity of the $9d \to 6p$ ($\lambda = 593.4$ nm) line. The $J_{9d}(t)$ dependences for a number of experimental conditions are shown in Fig. 3b, while the three-particle recombination coefficients $\beta_{e,He}(T_e)$ obtained by the described procedure are shown in Fig. 4b. It is seen from Fig. 4b that as the electron temperature is raised the decrease of the recombination coefficient $\beta_{e,He}(T_e)$ in a helium–xenon plasma differs substantially from the usual relation $\beta_e(T_e) \propto T_e^{-9/2}$ [as do also the relative intensities $J_\lambda(T_e)$ of the emission of the recombination-excited spectral lines (see Fig. 4a)].

In a number of cases, the measurements of the recombination coefficient $\beta_{e,He}$ on the basis of spectroscopic observations were repeated by determining $\beta_{e,He}$ directly from the decrease of the electron density $N_e(t)$ with time. This decrease was determined from the decrease of the current amplitude $i(t)$, see Eq. (26) [by application of a short rectangular electric-field pulse $E(t)$], at various instants of time t in the afterglow [18]. The recombination coefficients $\beta_{e,\mathrm{He}}$ obtained by the two methods were found to be close.

Note also that in the investigated range of pressures $P_{\mathrm{Xe}} = 0.01$–0.3 Torr no dependences of the radiation intensities on the xenon pressure were observed for the Xe* atom lines (for transitions from states with energies $|\varepsilon_{nl}| < \mathscr{E}_0$). This indicates that three-particle recombination of electrons with atomic Xe^+ ions in the gas Xe proper (the mechanism investigated in Ref. 4) is insignificant in an He/Xe mixture plasma under the considered experimental conditions ($N_{\mathrm{Xe}} \ll N_{\mathrm{He}}$).

5. CONCLUSION

It follows from the results of the calculations using Eqs. (23) and (24) that the experimental data of Sec. 4 on the relative radiation intensities of recombination-excited spectral lines of Xe atoms and on the coefficients of recombination of electrons with atomic Xe^+ ions in an He/Xe mixture plasma can be quantitatively accounted for by the recombination mechanism proposed in this paper (see Figs. 4a and 4b). In particular, under the experimental conditions ($\alpha = 6\cdot 10^{-8}$–$8\cdot 10^{-7}$ and $T_e = 400$–2500 K, see the caption of Fig. 4) we get from Eq. (22), for the characteristic energy ε_Y that demarcates two temperature regions with qualitatively different behavior of the recombination rates, the range $\varepsilon_Y = 700$–1000 K, which agrees with the measurement results. At low temperatures $T_e \lesssim 800$K there is realized the usual diffusion mechanism of three-particle electron-ion recombination with the free plasma electrons, and the recombination rates vary like $\Gamma \propto N_e^3 T_e^{-9/2}$ [see Eqs. (24) and (25a) and Figs. 4a and 4b]. At high temperatures $T_e \gtrsim 800$ K the decrease of the recombination coefficient with increase of T_e slows down noticeably, $\beta_{e,\mathrm{He}} \propto T_e^{-1.6}$, and the values of $\beta_{e,\mathrm{He}}$ turn out to be anomalously high. The "binary" coefficient $\alpha_{e,\mathrm{He}}$ $[\mathrm{cm}^3\cdot\mathrm{s}^{-1}] = \beta_{e,\mathrm{He}} N_e = \Gamma/N_e^2$ of the three-particle recombination depends weakly in this case on the electron density, and increases substantially when the concentration of the buffer gas is increased, $\alpha_{e,\mathrm{He}} \propto N_e^{0.3}\cdot N_{\mathrm{He}}^{0.7}$ [see Eqs. (24) and (25b) and Sec. 4]. We have thus demonstrated in this paper, experimentally and theoretically, the important role of collisions with neutral particles in electron-ion recombination of a noble-gas mixture plasma even at quite high degrees of ionization, up to $\alpha \approx 10^{-4}$.

Note that collisions with electrons and atoms of a buffer gas were simultaneously considered in the framework of the Fermi mechanism in a theoretical investigation of electron-ion recombination.[19] The temperature dependences of the recombination rates obtained there were substantially slower, $\alpha_{e,\mathrm{He}} \propto T_e^{-3/2}$, than in the case of only electron-electron collisions ($\alpha_e \propto T_e^{-9/2}$). In the model of Ref. 19, the plasma electrons assumed mainly the role of an ionizing factor that led to slowing down of the recombination flux produced by the neutral particles. This led to a decrease of the recombination coefficient $\alpha_{e,\mathrm{He}}$ $[\mathrm{cm}^3\cdot\mathrm{s}^{-1}] = \beta_{e,\mathrm{He}} N_e$ with increase of the electron density, in contrast to our experimental and theoretical result $\alpha_{e,\mathrm{He}} \propto N_e^{0.3}$. It must be emphasized in this connection that the presence of the here-investigated (see Sec. 2) mechanism of quenching of Rydberg atomic levels by neutral particles (a quenching more effective than by the traditional Fermi mechanism) leads to a qualitatively different character of electron-ion recombination, whereby the buffer-gas atoms enhance substantially the recombination flux formed by the free electrons of the plasma.

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[1)]We neglect in this paper the weak effect of the nonadiabatic coupling of the $|j\Omega,nl\rangle$ and $|j\Omega',n'l'\rangle$ states which are due to the Coriolis interaction of the rotations of the internuclear axis with the total angular momentum of the electron shell of the $\mathrm{YX}^+ + e^-$ quasimolecule. Estimates show that this interaction (see Ref. 5, p. 400 of Russian original) is much weaker than the here-considered Coulomb interaction V of the external and internal electrons.

[2)]Note that for transitions from nl levels with energies ε_{nl} exceeding the energy $\mathscr{E}_0$ of the ground vibrational level of the Xe_2^+ ion ($v = 0$) all the investigated time and temperature dependences of the Xe*-atom line intensities [$J_\lambda(t)$ and $J_\lambda(T_e)$] had the same form. This gives grounds for assuming that the intensities J are proportional to the recombination flux $\Gamma = \beta_{e,\mathrm{He}} N_e^2[\mathrm{Xe}^+]$ of electrons with atomic Xe^+ ions [see also (30)].

[1]A. V. Gurevich and L. P. Pitaevskiĭ, Zh. Eksp. Teor. Fiz. **46**, 1281 (1964) [Sov. Phys. JETP. **19**, 870 (1964)].

[2]L. P. Pitaevskiĭ, *ibid.* **42**, 1326 (1962) [**15**, 919 (1962)].

[3]L. M. Biberman, V. S. Vorob'ev, and I. T. Yakubov, *Kinetics of Nonequilibrium Low-Temperature Plasmas*, Consultants Bureau, 1987.

[4]V. S. Marchenko, Khim. Fizika **4**, 595 (1985).

[5]L. D. Landau and E. M. Lifshitz, *Quantum Mechanics, Nonrelativistic Theory*, Pergamon, 1977.

[6]D. Hausamann and H. Morgner, Mol. Phys. **54**, 1085 (1985).

[7]B. Brunetti, F. Vecchiocattivi, A. Anguilar-Navarro, and A. Solé, Chem. Phys. Lett. **126**, 245 (1986).

[8]I. I. Sobel'man, *Introduction to the Theory of Atomic Spectra*, Pergamon, 1973.

[9]H. A. Bethe and E. E. Salpeter, *Quantum Mechanics of One- and Two-Electron Systems*, Springer, 1958.

[10]E. Clementi and C. Roetti, Atom. Data and Nucl. Data Tables, Vol. 14, p. 177.

[11]M. Abramowitz and I. A. Stegun, *Handbook of Mathematical Functions*, Dover, 1964.

[12]V. S. Lebedev and V. S. Marchenko, Zh. Eksp. Teor. Fiz. **88**, 754 (1985) [Sov. Phys. JETP **61**, 443 (1985)].

[13]I. L. Beĭgman, *ibid.* **73**, 1729 (1977) [**46**, 908 (1977)].

[14]L. A. Vainshtein, I. I. Sobel'man, and E. A. Yukov, *Excitation of Atoms and Broadening of Spectral Lines*, Springer, 1981.

[15]V. A. Ivanov and Yu.E. Skoblo, Zh. Tekh. Fiz. **51**, 1386 (1981) [Sov. Phys. Tech. Phys. **26**, 796 (1981)].

[16]C. L. Chen, Phys. Rev. **131**, 2550 (1963).

[17]Y. J. Shiu, M. A. Biondi, and D. P. Sipler, Phys. Rev. **A15**, 494 (1977).

[18]V. A. Ivanov, Yu. E. Skoblo, and V. S. Sukhomlinov, Fiz. Plazmy **10**, 619 (1984) [Sov. J. Plasma Phys. **10**, 361 (1984)].

[19]C. B. Collins, Phys. Rev. **177**, 254 (1969).

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