The influence of quantum effects on inelastic ion–ion collisional excitations in dense, high-temperature plasmas

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\textit{New Journal of Physics} 7 (2005) 56

Received 19 August 2004
Published 15 February 2005
Online at http://www.njp.org/
doi:10.1088/1367-2630/7/1/056

Abstract. The ion–ion collision processes in dense, high-temperature plasmas are investigated. We focus upon the antiscreening channels of the excitation of a one-electron target by a one-electron projectile. In a pseudopotential framework we take into account plasma-screening effects as well as quantum effects of diffraction and symmetry. The semiclassical straight-line trajectory method is applied to the ion projectile path in order to calculate transition probabilities and cross sections. It is a well-known fact that screening reduces the transition probabilities of the ‘bare’ collision process. In contrast to this, the quantum effects incorporated here lead to an enhancement of the transition probabilities. In this paper, we study how these effects compete.

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1. Introduction

In recent years, atomic processes [1, 2] in dense high-temperature plasmas, such as ion–ion or electron–ion collisions, have received much attention [3]–[8]. These kinds of atomic phenomena can be a useful tool for plasma diagnostics which provide detailed information about plasma parameters. In plasma heating processes, the plasma heating method is given by directly injecting high power, energetic particles which then depose their energy onto the plasma. The principle of these processes is largely based on the above-mentioned collision processes, most importantly the ion–ion collision processes, and their cross-section data [8]. In astrophysical plasmas of compact objects and laboratory inertial confinement fusion plasmas, the range of the electron density and temperatures are known to be around \( n_e = 10^{20}–10^{23}\ \text{cm}^{-3} \) and \( T = 10^7–10^8\ \text{K} \), respectively. In this regime, the Debye length

\[
\Lambda = \sqrt{\frac{k_B T}{4\pi n_e e^2}},
\]

(1)

is greater than 10 times the first Bohr radius \( a_Z (= a_0/Z) \) of a hydrogenic ion with nuclear charge \( Z \) (\( k_B \) and \( e \) denote the Boltzmann constant and the electron charge, respectively). These plasmas can be classified as weakly coupled plasmas since the corresponding plasma coupling parameter \( \Gamma \) is smaller than unity. In recent years, collision processes in weakly coupled plasmas have been extensively studied using the well-known Debye–Hückel potential model [4, 5, 9] for coupling with plasma-screening effects. Also, the study on physical properties of dense high-temperature plasmas has been of great interest [10]. This is a result of common interest in investigations of astrophysical plasmas such as the interiors of the compact objects such as neutron stars and white dwarfs, which are the products of the final stages of stellar evolution. However, in high-density plasmas, the ratio between the average interparticle distance and the de Broglie thermal wavelength, governing the quantum nature, is small. Therefore, as has been pointed out, e.g. by Arkhipov et al [11, 12], in addition to plasma screening which governs the scattering at long distances, quantum effects of diffraction and symmetry at short distances should be included in a more correct treatment of such processes [13]–[15]. It is quite obvious that the properties of matter existing under such dense plasmas differ radically from the properties of a classical plasma. In these dense high-temperature plasmas, the interaction potential is different from the Debye–Hückel type because of strong collective effects and quantum mechanical effects of diffraction and symmetry on a level with plasma polarization effects [11, 16, 17]. Thus, in this...
context a pseudo-interaction potential for the particles in the plasma combining screening effects and quantum effects of diffraction and symmetry. Hence, in this paper we investigate ion–ion collisions in dense, high-temperature plasmas using this pseudopotential approach.

Irrespective of the effective interaction potential, the excitation of a one-electron target by a one-electron projectile is governed by two different mechanisms [1, 2]. These are directly related to the state of the projectile system after the collision and are called screening and antiscreening channels. In screening channels, the projectile electron remains in its ground state after the collision (elastic scattering of the projectile). However, in antiscreening channels the projectile electron is excited due to target-nucleus–projectile-electron interaction (inelastic scattering of the projectile). In this paper, we restrict ourselves to the more important antiscreening channels.

To simplify matters, the semiclassical straight-line trajectory method is applied [1, 2, 18] in order to calculate transition amplitudes and cross sections. Here, the motion of the projectile ion is described as a function of the classical impact parameter and the kinetic energy of the projectile ion.

The paper is organized as follows. In section 2, we derive the transition amplitude for ion–ion collisions in dense plasmas taking into consideration both plasma-screening and quantum effects. In section 3, we explicitly obtain the atomic form factors, transition amplitudes, and cross sections for the antiscreening channels. In section 4, we present the results for different plasma parameters. Section 5 is used for discussion and conclusions.

2. Transition amplitudes

In the semiclassical approximation, the excitation cross section in first Born approximation from an unperturbed atomic state \(|i⟩\equiv|\mathbf{i}_T, \mathbf{i}_P⟩\) to an excited state \(|f⟩\equiv|\mathbf{f}_T, \mathbf{f}_P⟩\) is given by [19]

$$\sigma = 2\pi \int_0^\infty b \, db |T_{\mathbf{f}_P, \mathbf{i}_P}^{\mathbf{f}_T, \mathbf{i}_T}(b)|^2,$$

where \(T_{\mathbf{f}_P, \mathbf{i}_P}^{\mathbf{f}_T, \mathbf{i}_T}(b)\) is the transition amplitude and \(b\) is the impact parameter. In the atomic wave functions introduced above, \(T\) and \(P\) stand for the target and projectile systems. The transition amplitude \(T_{\mathbf{f}_P, \mathbf{i}_P}^{\mathbf{f}_T, \mathbf{i}_T}(b)\) is given by the interaction potential \(V_{\text{int}}(\mathbf{r}_P, \mathbf{r}_T, \mathbf{R})\)

$$T_{\mathbf{f}_P, \mathbf{i}_P}^{\mathbf{f}_T, \mathbf{i}_T}(b) = -\frac{i}{\hbar} \int_{-\infty}^{\infty} dt \, e^{i\omega_{fi} t} \langle \mathbf{f}_T, \mathbf{f}_P|V_{\text{int}}(\mathbf{r}_P, \mathbf{r}_T, \mathbf{R})|\mathbf{i}_T, \mathbf{i}_P⟩,$$

where \(\omega_{fi} \equiv \Delta E/\hbar\) and \(\Delta E = (E_\mathbf{f}^p - E_\mathbf{i}^p + E_\mathbf{T}^T - E_\mathbf{T}^i)\) is the total energy change of the collision system, \(\mathbf{f}\) and \(\mathbf{i}\) denote the final and initial state, respectively. Figure 1 gives the coordinates applying to the present collision system. This semiclassical impact parameter method has strong appeal in aiding physical intuition since calculations based on this method are mathematically more tractable than the fully quantum mechanical treatments. Recently, an integro-differential equation for the effective potential of the particle interaction of semiclassical plasma, taking into account both quantum effects of diffraction, symmetry and plasma screening effects, was obtained on the basis of a sequential solution in semiclassical plasmas. Then, Baimbetov et al [13]–[15] proposed the use of the effective potential at short distances and the screened potential, treating three-particle correlation at large distances. Therefore, here we are using the effective projectile–target interaction potential that takes into account plasma-screening effects.
Figure 1. Schematic drawing of the collision system. The projectile nucleus \((Z_P)\) moves with a constant velocity \(v\) at an impact parameter \(b\). Here, \(R\) is the position of the projectile nucleus with respect to the target nucleus, the projectile electron has the coordinate \(r_P\) with respect to the projectile nucleus \(Z_P\) and the target electron has the coordinate \(r_T\) with respect to the target nucleus \(Z_T\).

at large distances as well as quantum effects of diffraction and symmetry at short distances. It reads [11, 12]

\[
V_{\text{int}}(r_P, r_T, R) = -\frac{Z_P Z_T e^2}{R} (e^{-R/\Lambda} - e^{-R/\lambda_{ii}}) - \frac{Z_T e^2}{|R + r_P|} (e^{-|R + r_P|/\Lambda} - e^{-|R + r_P|/\lambda_{ii}})
\]

\[
- \frac{Z_P e^2}{|R - r_T|} (e^{-|R - r_T|/\Lambda} - e^{-|R - r_T|/\lambda_{ii}}) + \frac{e^2}{|R + r_P - r_T|} (e^{-|R + r_P - r_T|/\Lambda} - e^{-|R + r_P - r_T|/\lambda_{ii}})
\]

\[
+ k_B T \ln 2 \exp \left[ \frac{-|R + r_P - r_T|^2}{\lambda_{ee}^2 \pi \ln 2} \right], \tag{4}
\]

where \(\Lambda\) is the Debye length (defined by equation (1)), \(\lambda_{ab} = \hbar / (2\pi \mu_{ab} k_B T)^{1/2}\) is the thermal de Broglie wavelength, \(\mu_{ab} = m_a m_b / (m_a + m_b)\) is the reduced mass of an a-b pair, respectively [\(a, b = e(\text{lectron}), i(\text{on})\)], \(k_B\) denotes the Boltzmann constant and \(T\) is the plasma temperature. Here and in the following, we assume that the projectile and target ion are of the same species, i.e. the charge numbers are the same \((Z_P = Z_T = Z)\). Note, equation (4) is appropriate for describing weakly coupled plasmas where additionally \(\lambda_{ii}, \lambda_{ei}, \lambda_{ee} \ll \Lambda\) holds [11]. If the terms including any \(\lambda_{ab}\) are neglected, i.e. ignoring the quantum effects, one attains the well-known static nonspherical Debye–Hückel interaction potential. If the collision velocity is fast with respect to the Bohr velocities of target and projectile electrons, the electron exchange effect can be neglected and the initial and final state wave functions can be written as products of target and projectile wave functions \(|f_T, i_P\rangle = \psi_{f_T}(r_T) \psi_{i_P}(r_P)\) and \(|i_T, f_P\rangle = \psi_{i_T}(r_T) \psi_{f_P}(r_P)\) [1].

Note that in our semiclassical treatment the position \(R\) of the projectile nucleus with respect to the target nucleus is treated classically. For target inelastic collisions \((f_T \neq i_T)\), this implies that the \(\frac{(Z_P)^2}{R} [e^{-R/\Lambda} - e^{-R/\lambda_{ii}}]\) and \(-\frac{Z_T^2}{|R + r_P|} [e^{-|R + r_P|/\Lambda} - e^{-|R + r_P|/\lambda_{ii}}]\) terms in equation (4) do not
contribute to the transition matrix due to the orthogonality of the initial and final states of the target system. Here, the transition amplitude becomes

$$T_{fi} = -\frac{i}{\hbar} \int_{-\infty}^{\infty} dt \ e^{i\omega_f t} \left[ -Ze^2 \delta_{fp,ip} \left( f_T \left| \frac{1}{|R - r_T|} \left( e^{-|R-r_T|/\lambda_e} - e^{-|R-r_T|/\lambda_w} \right) \right| i_T \right) ight. \\
+ e^2 \left( \left| f_T, f_P \right| \frac{1}{|R + r_p - r_T|} \left( e^{-|R+r_p-r_T|/\lambda_e} - e^{-|R+r_p-r_T|/\lambda_w} \right) \right| i_T, i_P \right) \\
+ \left( f_T, f_P \right| k_B T \ln 2 e^{-|R+r_p-r_T|^2/(\lambda_e^2 \pi \ln 2)} \right| i_T, i_P \right].$$

(5)

We now split the transition amplitude into two parts, namely, a part $T_{fi}^{DH}$ which is identical to the transition amplitude in the Debye–Hückel theory and a part $T_{fi}^{QM}$ comprising the quantum effects,

$$T_{fi} = T_{fi}^{DH} + T_{fi}^{QM}.$$  

(6)

In momentum space representation, these terms read

$$T_{fi}^{DH} = -\frac{ie^2}{2\pi^2 \hbar} \int_{-\infty}^{\infty} dt e^{i\omega_f t} \int d^3 q \frac{e^{-iqR}}{q^2 + \Lambda^{-2}} \left[ -\delta_{fp,ip} Z + G_{fp,ip}(-q) \right] F_{fi}^{on}(q)$$

(7)

and

$$T_{fi}^{QM} = -\frac{ie^2}{2\pi^2 \hbar} \int_{-\infty}^{\infty} dt e^{i\omega_f t} \int d^3 q e^{-iqR} \left[ \delta_{fp,ip} \frac{Z}{q^2 + \lambda_{ie}^{-2}} - \frac{1}{q^2 + \lambda_{ee}^{-2}} G_{fp,ip}(-q) \right] F_{fi}^{on}(q),$$

(8)

where $q$ is the momentum transfer and $G_{fp,ip}$ and $F_{fi}^{on}$ are the atomic form factors of the projectile and target systems, respectively,

$$G_{fp,ip}(-q) = \left( f_P \mid e^{-iq \cdot r_P} \mid i_P \right);$$

(9)

$$F_{fi}^{on}(q) = \left( f_T \mid e^{iq \cdot r_T} \mid i_T \right).$$

(10)

In the semiclassical straight-line (SL) trajectory method, we assume that the projectile is moving (classically) on a straight path given by

$$R(t) = b\hat{y} + vt\hat{z},$$

(11)

where $b$ is the impact parameter and $v$ is the collision velocity. For fast and heavy projectiles, deviations from a SL trajectory are negligible. For these high energy projectiles, the Massey parameter $(Ze)^2 \hbar v^{-1}$ [2] can be smaller than unity so that the first-order perturbation theory, which we use in this paper, is valid. The integration over time in equations (7) and (8) may now
be carried out,

\[
\int_{-\infty}^{\infty} dt \, e^{i\omega f t} e^{-i\mathbf{q} \cdot \mathbf{R}(t)} = e^{-i\mathbf{q} \cdot \mathbf{b}} 2\pi \delta(\mathbf{q} \cdot \mathbf{v} - \Delta). \tag{12}
\]

We proceed by setting \( \mathbf{q} = \mathbf{q}_\perp + \mathbf{q}_\parallel \hat{z} \), where \( \mathbf{q}_\perp \) and \( \mathbf{q}_\parallel \) are the perpendicular and parallel components of the momentum transfer with respect to the moving direction \( \hat{v} \). Equation (12) implies that \( \mathbf{q}_\parallel = \Delta / \hbar v \). This finally gives

\[
T_{fi}^{\text{DH}} = \frac{2ie^2}{\hbar v} \int_{0}^{\infty} q_\perp dq_\perp J_0(q_\perp b) F_{f_T,i_T}(q_\perp, \Delta / \hbar v) \left[ \delta_{f_T,i_T} Z - G_{f_T,i_T}(-q_\perp, -\Delta / \hbar v) \right] \tag{13}
\]

and

\[
T_{fi}^{\text{QM}} = -\frac{2ie^2}{\hbar v} \int_{0}^{\infty} q_\perp dq_\perp J_0(q_\perp b) F_{f_T,i_T}(q_\perp, \Delta / \hbar v) G_{f_T,i_T}(-q_\perp, -\Delta / \hbar v) \times \left[ \frac{Z\delta_{f_T,i_T}}{q_\perp^2 + (\Delta / \hbar v)^2 + \lambda_{ie}^{-2}} - \frac{1}{q_\perp^2 + (\Delta / \hbar v)^2 + \lambda_{ie}^{-2}} \right. \right. \\
\left. \left. + \frac{k_B T (\ln 2)^{5/2}}{4e^2} \pi^2 \lambda_{ie}^{3/2} e^{-(1/4)\lambda_{ie}^{-1} \ln(2)[q_\perp^2 + (\Delta / \hbar v)^2]} \right], \tag{14}
\]

where \( J_0 \) is the zeroth order Bessel function [20],

\[
J_0(x) = \frac{1}{2\pi} \int_{0}^{2\pi} d\phi \, e^{ix\cos\phi}. \tag{15}
\]

### 3. Antiscreening excitations

In the last section, we restricted our derivation of the transition amplitude to target inelastic scattering processes \( f_T \neq i_T \). In addition, we now concentrate on the antiscreening channels, i.e. we set \( f_T \neq i_T \) in the following (total inelastic processes). Consequently, the terms proportional to \( \delta_{f_T,i_T} \) in equations (13) and (14) drop out.

We now calculate the atomic form factors \( F_{f_T,i_T} \) and \( G_{f_T,i_T} \). As atomic wave functions we use unscreened hydrogenic wave functions since the main purpose of this paper is to investigate the inclusion of the quantum effects in the interaction potential [cf (4)]. The normalized radial parts of the lowest hydrogenic wave functions are [21]

\[
R_{1s}(r) = 2a_Z^{-3/2} e^{-r/2a_Z}, \tag{16}
\]

\[
R_{2s}(r) = \frac{1}{\sqrt{2}} a_Z^{-3/2} \left( 1 - \frac{r}{2a_Z} \right) e^{-r/2a_Z}, \tag{17}
\]

\[
R_{2p}(r) = \frac{1}{2\sqrt{2}} a_Z^{-5/2} r e^{-r/2a_Z}. \tag{18}
\]
In the following, we explicitly focus on the following three channels:

- the quadrupole–quadrupole (QQ) channel, i.e. $1s \rightarrow 2s$ (target) and $1s \rightarrow 2s$ (projectile) excitations,
- the dipole–quadrupole (DQ) channel, i.e. $1s \rightarrow 2p_0$ (target) and $1s \rightarrow 2s$ (projectile) excitations,
- the dipole–dipole (DD) channel, i.e. $1s \rightarrow 2p_0$ (target) and $1s \rightarrow 2p_0$ (projectile) excitations.

Inserting equations (16)–(18) into equations (9) and (10) we get the target and projectile form factors needed for calculating the transition amplitudes of the above-stated channels

$$G_{2s,1s}(-\bar{q}) = F_{2s,1s}(\bar{q}) = \frac{4\sqrt{2}}{a_0^2} \left[ q_{\perp}^2 + \frac{(\Delta/h\nu)^2}{q_{\perp}^2 + (\Delta/h\nu)^2 + (3/2a_0Z)^2} \right]^{3/2},$$

$$G_{2p_0,1s}(-\bar{q}) = -F_{2p_0,1s}(\bar{q}) = -\frac{i6\sqrt{2}}{a_0^2} \left[ q_{\perp}^2 + \frac{(\Delta/h\nu)^2}{q_{\perp}^2 + (\Delta/h\nu)^2 + (3/2a_0Z)^2} \right]^{1/2}.$$

Finally, the transition amplitudes for the QQ, DQ, and DD channels are given by

$$T_{QQ}(\bar{\epsilon}, \bar{b}) = \frac{2^8}{3Z} \sqrt{\frac{\alpha}{\bar{\epsilon}}} \int_{0}^{\infty} \tilde{d}q_{\perp} J_0(q_{\perp} \bar{b}) q_{\perp} [\frac{q_{\perp}^2 + (\alpha/\bar{\epsilon})^2}{q_{\perp}^2 + (\alpha/\bar{\epsilon}) + (9/4)}]^{3/2} K_{\bar{\epsilon}}(q_{\perp}),$$

$$T_{DQ}(\bar{\epsilon}, \bar{b}) = \frac{i2^7}{Z} \sqrt{\frac{\alpha}{\bar{\epsilon}}} \int_{0}^{\infty} \tilde{d}q_{\perp} J_0(q_{\perp} \bar{b}) q_{\perp} [\frac{q_{\perp}^2 + (\alpha/\bar{\epsilon})^2}{q_{\perp}^2 + (\alpha/\bar{\epsilon}) + (9/4)}] \bar{\epsilon} K_{\bar{\epsilon}}(q_{\perp}),$$

$$T_{DD}(\bar{\epsilon}, \bar{b}) = \frac{2^6 3}{Z} \sqrt{\frac{\alpha}{\bar{\epsilon}}} \int_{0}^{\infty} \tilde{d}q_{\perp} J_0(q_{\perp} \bar{b}) q_{\perp} [\frac{q_{\perp}^2 + (\alpha/\bar{\epsilon})^2}{q_{\perp}^2 + (\alpha/\bar{\epsilon}) + (9/4)}] K_{\bar{\epsilon}}(q_{\perp}),$$

where

$$K_{\bar{\epsilon}}(q_{\perp}) = -\frac{1}{q_{\perp}^2 + (\alpha/\bar{\epsilon}) + \Lambda^{-2}} + \frac{1}{q_{\perp}^2 + (\alpha/\bar{\epsilon}) + \bar{\lambda}_{ee}^{-2}} - \bar{\lambda}_{ee} \pi \ln(2)^{5/2} \exp\{-(1/4)\bar{\lambda}_{ee}^2 \pi \ln(2)[q_{\perp}^2 + (\alpha/\bar{\epsilon})]\}.$$  

Here, we use scaled quantities, i.e., lengths and energies are measured in units of Bohr radii, $a_0 = a_0/Z$, and Rydberg’s, $Z^2$Ry, respectively. The scaled quantities are indicated by a bar; $\bar{b} = b/a_0$, $\bar{q}_{\perp} = q_{\perp} a_0$, etc. Besides this, we introduced the scaled collision energy $\bar{\epsilon}$ and the constant $\alpha$ which are given by $\bar{\epsilon} = \mu v^2/(2Z^2$Ry) and $\alpha = 9\mu/16m$, where $\mu$ is the reduced mass of the collision system, and $m$ is the electron mass. The scaled cross section reads (cf equation (2))

$$\bar{\sigma}_{\text{ch}}(\bar{\epsilon}) = \sigma_{\text{ch}}(\bar{\epsilon})/a_0^2 = 2\pi \int_{0}^{\infty} \tilde{b} \, d\tilde{b} |T_{\text{ch}}(\bar{\epsilon}, \bar{b})|^2,$$

where the superscript ‘ch’ stands for either QQ, DQ or DD channel.
Figure 2. (a) The QQ channel cross section $\bar{\sigma}^{QQ}(\bar{\varepsilon})$ as a function of the scaled collision energy $\bar{\varepsilon}$; (b) the difference $\bar{\sigma}^{QQ}(\bar{\varepsilon}) - \bar{\sigma}^{QQ}_{0}(\bar{\varepsilon})$, where $\bar{\sigma}^{QQ}_{0} = \bar{\sigma}^{QQ}[\bar{\Lambda} = \bar{\lambda}_{ee} = 0]$; $\bar{\Lambda} = 0$ and $\bar{\lambda}_{ee} = 0$; $\bar{\Lambda} = 10$ and $\bar{\lambda}_{ee} = 0$; $\cdots$, $\bar{\Lambda} = 0$ and $\bar{\lambda}_{ee} = 0.1003$; $\bar{\Lambda} = 10$ and $\bar{\lambda}_{ee} = 0.0317$; $\bar{\Lambda} = 10$ and $\bar{\lambda}_{ee} = 0.1003$.

Figure 3. (a) The DQ channel cross section $\bar{\sigma}^{DQ}(\bar{\varepsilon})$ as a function of the scaled collision energy $\bar{\varepsilon}$; (b) the difference $\bar{\sigma}^{DQ}(\bar{\varepsilon}) - \bar{\sigma}^{DQ}_{0}(\bar{\varepsilon})$, where $\bar{\sigma}^{DQ}_{0} = \bar{\sigma}^{DQ}[\bar{\Lambda} = \bar{\lambda}_{ee} = 0]$; $\bar{\Lambda} = 0$ and $\bar{\lambda}_{ee} = 0$; $\bar{\Lambda} = 10$ and $\bar{\lambda}_{ee} = 0$; $\cdots$, $\bar{\Lambda} = 0$ and $\bar{\lambda}_{ee} = 0.1003$; $\bar{\Lambda} = 10$ and $\bar{\lambda}_{ee} = 0.0317$; $\bar{\Lambda} = 10$ and $\bar{\lambda}_{ee} = 0.1003$. 
4. Plasma screening versus quantum effects

In our calculations, the scaled Debye length $\tilde{\Lambda}$ is set to $\tilde{\Lambda} = 10$. For the de Broglie wavelength $\tilde{\lambda}_{ee}$ we consider two cases, namely, $\tilde{\lambda}_{ee} = 0.0317$ and $\tilde{\lambda}_{ee} = 0.1003$ which refer to $T = 10^8$ and $10^7$ K, respectively. Note that a lowering of the temperature from $T = 10^8$ and $10^7$ K, while keeping the Debye length fixed implies a decrease of the particle density in the plasma. For comparison, the cross sections $\tilde{\sigma}^{ch}_0$ for the ‘bare’ processes (i.e. excluding plasma screening as well as quantum effects) and the cross sections excluding one of either screening or quantum effects, are also calculated. For a detailed discussion of the influence of plasma screening effects on totally inelastic ion–ion scattering, we refer to a recent paper by Yoon and Jung [5]. Solely taking into account plasma screening (Debye–Hückel model) leads to a lowering of the bare cross sections.

In figures 2–4, we present the different cross sections for the above-stated plasma parameters. In the temperature regime we investigated, the quantum effects always enhance the bare cross sections, therefore, a competition between quantum effects and plasma polarization effects (plasma screening) takes place. For high enough collision energies, plasma screening is dominating in all the cases and hence the bare cross sections are reduced in the high energy regime. However, going to lower energies we find that the quantum effects provide...
Figure 5. Cross section $\bar{\sigma}(\bar{\epsilon})$ as a function of the scaled energy $\bar{\epsilon}$ for plasma parameters $\bar{\Lambda} = 10$ and $\bar{\lambda}_{ee} = 0.1003$. The solid line represents the QQ channel, the dashed line the DQ channel, and the dotted line the DD channel.

for an overall enhancement of the bare cross sections. For $\bar{\lambda}_{ee} = 0.0317$ ($T = 10^8$ K) quantum effects are less important than for $\bar{\lambda}_{ee} = 0.1003$ ($T = 10^8$ K). Besides these general aspects, we find some quantitative differences between the individual channels we were looking at. At first, figure 5 shows the difference in magnitude between the QQ, DQ, and QQ channels for $(\bar{\Lambda} = 10, \bar{\lambda}_{ee} = 0.1003)$. It is clear from this that the DD channel by far dominates. Compared to the other channels, in the QQ channel the quantum effects are most important whereas in the DD channel they play the least role.

5. Summary

In this paper, we investigated totally inelastic ion–ion scattering processes in dense, high-temperature plasmas. More precisely, we studied the excitations of a one-electron target ion by a one-electron projectile ion. In an effective projectile–target interaction potential framework [11, 12] we took into account plasma polarization effects, i.e. plasma screening, as well as quantum effects of diffraction and symmetry. We explicitly calculated the cross sections for three different channels, namely, the quadrupole–quadrupole (QQ), the dipole–quadrupole (DQ), and the dipole–dipole (DD) excitation channels. We used a semiclassical straight-line trajectory method for the projectile ion. We find that the quantum effects are of the same order as the plasma polarization effects. However, in contrast to the latter, they lead (at least in the temperature regime we focused on) to an enhancement of the bare cross sections, showing an antiscreening effect. This in turn leads to a competition where plasma polarization dominates for higher collision energies whereas for lower energies the quantum effects outweigh. Moreover, we find that in the QQ channel the quantum effects have the strongest influence on the cross section compared to the other two channels.
These results provide a general description of the inelastic ion–ion collisional excitations in dense, high-temperature plasma.

Acknowledgments

This work was supported by the Korea Institute of Science and Technology Information (KISTI) through the construction of atomic database for the industrial application plasma project.

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