Electron transport in dense degenerate plasmas

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Abstract. Within a unified approach, a method for calculating the tensors of electrical conductivity, the Seebeck coefficient and thermal conductivity of a nonideal plasma in a magnetic field were considered. Under this unified approach the kinetic coefficients are calculated together with the equation of state for a nonideal plasma within the framework of a quasi-chemical model. Various methods for determining the Coulomb logarithm in the kinetic theory of transport and various options for choosing the boundary value of the wave number of electrons are considered. The scattering of electrons by ions using the phase shift method has been considered and the appearance of values of the Coulomb logarithm less than unity are demonstrated. Electron scattering by the phase shift method is considered using the Buckingham potential which permits to describe the Ramsauer minimum in the transport cross section for electron scattering by noble gas atoms.

1. Introduction

In [1], the thermo-electrophysical properties of plasma at pressures of the megabar range without a magnetic field were studied. In [2], it was developed a method for calculating tensors of electrical conductivity, Seebeck coefficient (thermo-EMF) and thermal conductivity of nonideal plasma. In [3], various methods for determining the Coulomb logarithm in kinetic transport theory were considered. The correlation of ions was taken into account on the basis of the numerical solution of the Ornstein-Zernike (OZ) integral equation in the hypernetted chain (HNC) approximation.

The study of the Coulomb logarithm in the kinetic theory of electron transport have been continued in [4]. In this paper, a brief description will be made of the main results of papers [1–4], and the main attention will be paid to the quantum mechanical calculation of electron scattering by noble gas atoms using the phase shift method (PSM), which is important for calculating transport coefficients in a dense degenerate plasma, since the scattering of electrons occurs not only on ions, but also on more numerous atoms.

2. Coulomb logarithm in nonideal and degenerate plasma

In plasma physics, the long-range nature of the Coulomb interaction leads to the appearance of various kinds of divergent integrals. For example, this leads to a divergence of the transport cross sections for collisions of electrons and ions with charged particles, and the divergence occurs both toward small values of the impact parameter and toward large values (see [5]).
In [7], S.V. Temko derived the Fokker-Planck equation for plasma based on the Bogolyubov chain of equations and the following expression for the Coulomb logarithm he obtained

\[
\Lambda = \frac{1}{2} \ln \left( 1 + \frac{R_D^2}{r_0^2} \right),
\]

(1)

which, in contrast to the expression obtained in [5], does not give negative values of the Coulomb logarithm. Here \( R_D \) is the largest impact parameter equal to the Debye screening radius:

\[
R_D = k_D^{-1}, \quad k_D^2 = k_{De}^2 + k_{Di}^2, \quad k_{De}^2 = \frac{4\pi e^2 n_e}{T_e}, \quad k_{Di}^2 = \frac{4\pi z_i^2 e^2 n_i}{T_i},
\]

(2)

\( r_0 \) is the Landau radius defined as

\[
r_0 = e^2 z_i / T_e,
\]

(3)

e is the elementary charge, \( z_i \) is the charge number of the ion, \( T_e, T_i \) are the electron and ion temperatures in energetic units, \( n_e, n_i \) are the number densities of electrons and ions, respectively.

In [1], at determining the electrical conductivity of a dense hydrogen plasma, we also used an approach based on the Coulomb logarithm, while its values in the region of strong plasma ionization turned out to be noticeably less than unity, which raises certain questions about the applicability of the approach used to describe the transport electrons in this range of parameters. Therefore, in [3] a detailed consideration of this issue was carried out. It should be noted that scattering in a nonideal plasma is many-particle, and the accuracy of modern experiments does not allow one to determine even the exact order of magnitude of the Coulomb logarithm.

### 2.1. Transport cross section for scattering of electrons by ions

J. Zyman in [8] introduced and actively used the concept of an “atom”, an ion with a shielding electron cloud, to describe the transport properties of electrons in metals. In this case, electron scattering occurs on such an “atom” with a pseudopotential equal to the screened Debye one. And the interaction of the “atoms” with each other is described by the same potential. This theory was very successful in describing the transport properties of electrons in various metals and their alloys. Therefore, in [3], we apply the same approach.

Electron screening in plasma is generally defined by the expression (see [9,10])

\[
k_{De}^2 = \frac{8\pi e^2}{T_e k^3} \lambda^3 \mathcal{F}_{-1/2} (\eta_e),
\]

(4)

where \( \eta_e = \mu_e / T_e \), \( \mu_e \) is the chemical potential of the electron gas, \( \mathcal{F}_{-1/2} (\eta_e) \) is the Fermi-Dirac integral defined by the relation [11]:

\[
\mathcal{F}_k (\eta_e) = \frac{1}{\Gamma (k+1)} \int_0^\infty \frac{x^k dx}{e^{x-\eta_e} + 1}.
\]

(5)

\( \Gamma (x) \) is the gamma function. The electron number density is determined by the expression:

\[
n_e = \frac{2}{\lambda^3} \mathcal{F}_{1/2} (\eta_e),
\]

(6)

From (4) and (6) we find

\[
k_{De}^2 = k_{De,0}^2 \frac{\mathcal{F}_{-1/2} (\eta_e)}{\mathcal{F}_{1/2} (\eta_e)},
\]

(7)
where $k_{De,0}$ is the electron screening constant in the nondegenerate case (see (2))

$$k_{De,0}^2 = \frac{4\pi e^2 n_e}{T_e}.$$  \hfill (8)

Note that at room temperature and $n_e \leq 10^{18} \text{ cm}^{-3}$ the electron screening constant does not practically differ from (8), and in the strongly degenerate case

$$k_{De}^2 = \frac{m_e e^2}{\pi^2 \hbar^2} \left(\frac{3n_e}{\pi}\right)^{1/3}.$$  \hfill (9)

Note that Fermi-Dirac integrals were calculated according to [11].

The transport cross section for electron scattering by ions, taking into account the ion-ion correlation, is determined by the expression [12]:

$$Q_{ei} (\varepsilon_e) = \frac{\pi z_i^2 e^4}{\varepsilon_e^2} \Lambda_{ei},$$  \hfill (10)

where $\varepsilon_e$ is the electron energy, $\Lambda_{ei}$ is the Coulomb logarithm:

$$\Lambda_{ei} = \int_0^{q_m} \frac{k^3}{(k^2 + k_s^2)^2} S_i (k) \, dk,$$  \hfill (11)

$S_i (k)$ is the static structure factor (SSF) describing the ion-ion correlation with the interaction potential

$$U_{ii} (R) = \frac{\varepsilon_i^2 e^2}{R} e^{-k_s R},$$  \hfill (12)

$R$ is the distance between ions, $k_s$ is the screening constant, $q_m$ is the maximum value of the wavenumber (determined by the maximum value of the electron momentum transferred to the ion in the collision). In [8], it was supposed that $q_m = 2k_F$, where $k_F$ is the electron wavenumber at the Fermi surface: $k_F = (3\pi^2 n_e)^{1/3}$. In [1], the expression $q_m = \min \left\{ E_T / z_i e^2, 2/\lambda_e \right\}$ was used. Here in the nondegenerate case $E_T$ was assumed to be equal to the temperature (in energy units), and in the degenerate case to the Fermi energy. The de Broglie wavelength $\lambda_e$ in these cases was also determined using the temperature or the Fermi energy.

In a rarefied plasma, the ion-ion correlation can be neglected: $S_i \approx 1$. In this case from (11) we get (see [10,13]):

$$\Lambda_{ei} = \frac{1}{2} \left[ \ln (1 + \chi_i) - \frac{\chi_i}{1 + \chi_i} \right],$$  \hfill (13)

where $\chi_i = (q_m / k_s)^2$. With SSF in Debye approximation [14] from (11) it was found [1]:

$$\Lambda_{ei} = \frac{1}{2} \left[ \ln (1 + \chi_i) - \frac{\chi_i}{1 + \chi_i} - \frac{1}{2} \frac{\chi_i^2}{k_s^2} \frac{4\pi n_e e^2}{T} \frac{\lambda_i^2}{(1 + \chi_i)^2} \right].$$  \hfill (14)

To determine SSF of ions, we numerically solved the OZ equation [15–17]:

$$h (r) = C (r) + n_i \int h (r_1) C (|r - r_1|) \, dr_1,$$  \hfill (15)

where $g(r) = 1 + h(r)$ is the pair correlation function, $C(r)$ is the direct correlation function. To close the OZ equation, we used the HNC approximation [18]:

$$C (r) = \exp \left[ -\frac{U_{ii} (r)}{T} + \gamma (r) \right] - \gamma (r) - 1,$$  \hfill (16)
where $\gamma(r) = h(r) - C(r)$. In our case, the Debye potential (12) was used to describe the interaction of ions. The HNC approximation turns out to be accurate enough to describe Coulomb systems (see, for example, [19]).

The SSF is related to the pair correlation function by the Fourier transform [16]

$$S(k) = 1 + n_i \int [g(r) - 1] e^{-ikr} dr = 1 + \frac{4\pi n_i}{k} \int_0^\infty h(r) \sin(kr) r dr.$$  \hspace{1cm} (17)

The OZ equation was solved by the iteration method [20,21]. To accelerate the convergence, we used the procedure proposed in [22].

2.2. Numerical calculations of the Coulomb logarithm

In [3], two variants of the screening constant were considered: with allowance for ($k_s = k_D$) and without allowance for the ion contribution ($k_s = k_{De}$). The value of the boundary wave number $q_m$ was set either as in [1]:

$$q_m = q_{mS} \equiv \min \left( r_0^{-1} \sqrt{1 + (\varepsilon_F/T)^2}, 2k_E \right),$$  \hspace{1cm} (18)

or as in [8]:

$$q_m = q_{mZ} \equiv \sqrt{r_0^{-2} + 4k_E^2},$$  \hspace{1cm} (19)

where $k_E$ is the wave number determined from the mean kinetic energy of electrons:

$$k_E = \left(\frac{10}{3} \frac{\langle \varepsilon_e \rangle m_e}{\hbar^2}\right)^{1/2}, \quad \langle \varepsilon_e \rangle = \frac{3}{2} T \frac{F_{3/2}(\eta_e)}{F_{1/2}(\eta_e)}.$$  \hspace{1cm} (20)

In a non-degenerate plasma $\langle \varepsilon_e \rangle = \frac{3}{3} T$, and in a strongly degenerate one $\langle \varepsilon_e \rangle = \frac{3}{2} \varepsilon_F$, where $\varepsilon_F$ is the Fermi energy

$$\varepsilon_F = \left(3\pi^2 n_e\right)^{2/3} \frac{\hbar^2}{2m_e}.$$  \hspace{1cm} (21)

As a result, in [3], calculations were carried out for four options for choosing the maximum value of the wavenumber and the screening constant. But in the light of [23, 24], in which it is shown that the ion-ion correlation is determined by the Debye potential only with the electron screening constant, further we will consider only the options with $k_s = k_{De}$ with two different definitions of the boundary value of the wave number: $q_m = q_{mZ}$ and $q_m = q_{mS}$.

Figure 1 compares the two above-mentioned calculation options. It can be seen that in the calculations with $q_m = q_{mS}$ the Coulomb logarithm turn out to be noticeably smaller than in the calculations with $q_m = q_{mZ}$. As can be seen from figure 1, taking into account the ion-ion correlation turns out to be important in both cases. A similar picture was observed in calculations for $T = 300, 1500$ and $3000$ K.

Along with the Coulomb logarithm, in [3], calculations of the plasma electrical conductivity were carried out with four variants of determining the Coulomb logarithm. Specific electrical conductivity was determined within the framework of the Lorentz-Bloch model [1, 12]

$$\sigma_T = \frac{4e^2T^{-3/2}}{3\sqrt{\pi} m_e \lambda_e^3} \int_0^\infty \frac{\varepsilon_e^{3/2}}{v_m(\varepsilon_e)} \left[-\frac{\partial f_0}{\partial \varepsilon_e}\right] d\varepsilon_e,$$  \hspace{1cm} (22)
Figure 1. The Coulomb logarithm as a function of $n_e$ at $z_i = 1$, $k_s = k_{De}$ and $T = 5000$ K with $q_m = q_mZ$ (1) and $q_m = q_mS$ (2) calculated by numerical integration of (11) using the SSF from OZ and HNC equations (——), from (13) (— · —) and from (14) (······).

Figure 2. The conductivity of hydrogen plasma for $T = 5000$ K, $z_i = 1$, $k_s = k_{De}$ calculated with $q_m = q_mZ$ (1) and $q_m = q_mS$ (2), with $\Lambda_{ei} = 1$ (3); numerical data [1] (4), experimental data [25,26] (○) and [27] (△).

where $\nu_m$ is the transport frequency of electron collisions:

$$\nu_m (\varepsilon_e) = \sqrt{\frac{2\varepsilon_e}{m_e}} [n_i Q_{ei} (\varepsilon_e) + n_a Q_{ea} (\varepsilon_e)],$$

where $Q_{ea}$ and $Q_{ei}$ are the transport cross sections of electron scattering on atoms with the number density $n_a$ and ions with the number density $n_i$. If the plasma contains several types of atoms (molecules) and/or ions, then summation should be performed over them. In [3], only the case when collisions of electrons with singly ionized ions predominate was considered. In general, collisions with atoms and molecules should be taken into account, which will lead to a decrease in the electrical conductivity of the plasma. Therefore, the values presented in [3] determine only the upper limit of the electrical conductivity.

In the case of a predominance of electron-ion collisions, we have

$$\nu_m (\varepsilon_e) = \frac{\pi z_i^2 e^4 \Lambda_{ei}}{m_e} \sqrt{\frac{2}{m_e} n_i \varepsilon_e^{-3/2}}.$$  

Using (23), from (22) we get

$$\sigma_T = \frac{2T}{\pi z_i e^2 \Lambda_{ei}} \sqrt{\frac{8T}{\pi m_e}} F_2 (\eta_e) F_1 / 2 (\eta_e).$$  

The dependencies of the electrical conductivity for two options for choosing the boundary value of the wave number are shown in figure 2. It is set that a screening constant is equal to the electronic one. It can be seen that the smallest values of electrical conductivity occur in calculations with $q_m = q_mZ$, to which, at high electron number densities, the calculation data with $q_m = q_mS$ begin to approach. It can be seen from figure 2 that the electrical conductivity with a unit value of the Coulomb logarithm at $n_e > 10^{17}$ cm$^{-3}$ turn out to be less than in
calculations with the structure factor found in numerical solution of the OZ equation in the HNC approximation.

Figure 2 also shows the plasma conductivity data calculated in [1] taking into account the contribution of neutral atoms and molecules to electron scattering, and the Coulomb logarithm was calculated using (14). In this case, the contribution of neutral atoms and molecules to the total scattering cross section turned out to be overwhelming. Figure 2 also shows the experimental data measured in [25–29]. It can be seen that the calculated values of the conductivity are in fairly good agreement with the experimental data [25–29], the accuracy of which does not allow making the final choice in favor of one or another models for calculating the Coulomb logarithm.

3. Electrical conductivity of a nonideal plasma in a magnetic field

To calculate the electrical conductivity, thermo-EMF and thermal conductivity of a nonideal plasma in a magnetic field from moderate densities to the region of strong electron degeneracy in [2], a unified approach from [1] was used. The electrical conductivity tensor $\hat{\sigma}$ in the magnetic field $B$ directed along the axis $z$ has the form

$$\hat{\sigma} = \begin{pmatrix} \sigma_\perp & -\sigma_H & 0 \\ \sigma_H & \sigma_\perp & 0 \\ 0 & 0 & \sigma_\parallel \end{pmatrix},$$

(26)

where the tensor components within the $\tau$-approximation (Lorentz-Bloch model) [1,12] defined by expressions:

$$\sigma_\perp = A_\sigma \int_0^\infty \frac{\epsilon^{3/2}}{\Omega^2 + \nu_m^2(\epsilon)} \left( -\frac{\partial f_0}{\partial \epsilon} \right) d\epsilon, \quad \sigma_\parallel = A_\sigma \int_0^\infty \frac{\epsilon^{3/2}}{\nu_m(\epsilon)} \left( -\frac{\partial f_0}{\partial \epsilon} \right) d\epsilon,$$

$$\sigma_H = A_\sigma \int_0^\infty \frac{\epsilon^{3/2}\Omega}{\Omega^2 + \nu_m^2(\epsilon)} \left( -\frac{\partial f_0}{\partial \epsilon} \right) d\epsilon, \quad A_\sigma = \frac{4e^2T^{-3/2}}{3\sqrt{\pi}m_e \bar{\lambda}^3},$$

(27)

Here $f_0$ is the energy distribution function of electrons (the Fermi-Dirac distribution):

$$f_0(\epsilon) = \frac{1}{1 + e^{(\epsilon-\mu)/T}},$$

(28)

$\nu_m$ is the electron collision frequency defined above by (23), $\Omega = eB/cm_e$ is the Larmor frequency. As in [1,12], the transport cross sections for scattering on ions are calculated taking into account the ion-ion correlation.

The tensor $\hat{\tau}$, describing the electric current proportional to temperature gradient, has the same shape as $\hat{\sigma}$ with the components

$$\tau_\perp = k_B e A_\tau \int_0^\infty \frac{\epsilon^{3/2}\nu_m(\epsilon)}{\Omega^2 + \nu_m^2(\epsilon)} \left( \frac{\mu - \epsilon}{T} \right) \left( -\frac{\partial f_0}{\partial \epsilon} \right) d\epsilon, \quad \tau_\parallel = k_B e A_\tau \int_0^\infty \frac{\epsilon^{3/2}}{\nu_m(\epsilon)} \left( \frac{\mu - \epsilon}{T} \right) \left( -\frac{\partial f_0}{\partial \epsilon} \right) d\epsilon,$$

$$\tau_H = k_B e A_\tau \int_0^\infty \frac{\epsilon^{3/2}\Omega}{\Omega^2 + \nu_m^2(\epsilon)} \left( \frac{\mu - \epsilon}{T} \right) \left( -\frac{\partial f_0}{\partial \epsilon} \right) d\epsilon, \quad A_\tau = \frac{\sqrt{8m_e}}{3\pi^2 \hbar^3},$$

(29)

where $k_B$ is the Boltzmann constant. The Seebeck tensor $\hat{\alpha}$ is the product of the inverse electrical conductivity matrix and the matrix $\hat{\tau}$ (cf. [1,30]): $\hat{\alpha} = \hat{\sigma}^{-1}\hat{\tau}$. To describe the thermal
conductivity tensor of the plasma in the magnetic field, in [2] it was introduced the auxiliary tensor $\hat{\gamma}$ for the heat flux proportional to the temperature gradient in a given electric field, that has the same shape as $\hat{\sigma}$ with the components:

$$\gamma_\perp = \frac{A_\tau}{T} \int_0^\infty \frac{e^{3/2} \nu_m(\epsilon) (\mu - \epsilon)^2}{\Omega^2 + \nu_m^2(\epsilon)} \left( - \frac{\partial f_0}{\partial \epsilon} \right) d\epsilon$$

$$\gamma_\parallel = \frac{A_\tau}{T} \int_0^\infty \frac{\epsilon^{3/2} \Omega}{\nu_m(\epsilon) (\mu - \epsilon)^2} \left( - \frac{\partial f_0}{\partial \epsilon} \right) d\epsilon, \quad \gamma_H = \frac{A_\tau}{T} \int_0^\infty \frac{\epsilon^{3/2} \Omega}{\nu_m(\epsilon) (\mu - \epsilon)^2} \left( - \frac{\partial f_0}{\partial \epsilon} \right) d\epsilon, \quad (30)$$

The thermal conductivity tensor $\hat{\kappa}$ is expressed as $\hat{\kappa} = \hat{\gamma} - T \hat{\tau} \hat{\alpha}$. Various estimates of the tensor components $\hat{\sigma}$, $\hat{\tau}$ and $\hat{\kappa}$ in a magnetized plasma under different conditions are given in the papers [1,2].

Calculations have shown (for more details see [1]) that the values of thermo-EMF in hydrogen reach values of several thousand $\mu V/(K \cdot cm)$, i.e. exceed the typical values for semiconductors [31]. At densities below 18 g/cm$^3$, the xenon plasma is almost unmagnetized because of the prevailing scattering on neutral atoms. For this reason, the longitudinal electrical conductivity coincides with the transverse electrical conductivity, and both conductivities are much higher than the Hall component $\sigma_H$ of $\hat{\sigma}$. The observed maximum of the components of the electrical and thermal conductivities was due to the Ramsauer effect.

The influence of this effect is particularly strong, up to a change in sign, on the Seebeck coefficient (see figure 3, where symbols show the absolute values of the components of the tensor in the region of negative values). The effect of the medium at the indicated densities has not yet been taken into account. The calculations of the effect of the medium on the Ramsauer minimum in [32,33] refer to liquid xenon at densities near 3 g/cm$^3$. Since these calculations predict an increase in cross sections at energies above the Ramsauer minimum, a noticeable influence of this effect on galvano- and thermoelectric components in the nonideal degenerate xenon plasma can be expected. This prediction should be experimentally tested.

![Figure 3](image-url)

**Figure 3.** Components $\alpha_\parallel$ (1), $\alpha_H$ (2, 3), and $\alpha_\perp$ (4, 5) of the Seebeck tensor of the nonideal xenon plasma on the $T = 2000$ K isotherm in a magnetic field of $B = 10$ kG (2, 4) and 5 MG (3, 5); the absolute values of the negative Seebeck coefficients $|\alpha_\parallel|$ (●), $|\alpha_\perp|$ (△), and $|\alpha_H|$ (□) at $B = 5$ MG and $|\alpha_H|$ (■) at $B = 10$ kG ($\alpha_\perp$ values at 10 kG and 5 MG in the negative region almost coincide with each other, cf. also lines 4 and 5).

In the density range where the number of neutral atoms becomes negligibly small, electrons are strongly magnetized, which results in a large, many orders of magnitude, difference in the components of the tensors $\hat{\sigma}$, $\hat{\alpha}$ and $\hat{\kappa}$. This effect is due to a small value of the Coulomb logarithm in a strongly nonideal plasma [1] and should also be experimentally tested. It is noteworthy that the same result leads to anomalously high values of the longitudinal electrical and thermal conductivities, which are about five orders of magnitude higher than these values in copper and in TOKAMAK plasma, where the Coulomb logarithm is much larger than unity.
Figure 4. The total cross sections calculated by the phase shift method (1) and in the Born approximation (2), the PSM transport cross sections (3), the PSM total cross sections (4) and the PSM transport cross sections (5) for \( z_i = -1 \) as a function of electron energy at \( n_{e0} = n_{i0} = 10^{18} \) (a) and \( 10^{20} \) cm\(^{-3} \) (b).

4. Scattering of electrons by ions and atoms of noble gases

In [4], the behavior of the cross sections of electron scattering by ions was studied in detail based on the phase shift method (PSM) and the transport cross sections for scattering of electrons by ions was determined taking into account the correlation of ions on the basis of the integral OZ and HNC equations. Figure 4 shows the total and transport cross sections of electrons on positively and negatively (mainly for comparison) charged ions with an ionization charge state \( z_i \) with a Debye interaction potential with an electron screening constant. In figure 4, it can be seen that the dependences of both the total and transport cross sections for scattering by positive ions on the incident electron energy have extrema, which is associated with the influence of bound states [34]. For scattering by negative ions, as can be seen in figure 4, such features are not observed.

Figure 5 shows the values of the “Coulomb” logarithm,

\[
\Lambda(\varepsilon) = \frac{\sigma_{tr}\varepsilon^2}{\pi z_i^2 e^4},
\]

(31)
determined according to (10) from \( \sigma_{tr} \) calculated by PSM at \( S_i = 1 \). The horizontal lines (5) and (6) show the Coulomb logarithm at the specified values of the electron number densities, calculated for \( S_i = 1 \) from (13) at \( q_m = q_mZ \) (19) and at \( q_m = q_mS \) (18). These values are also shown in table 1. Comparison of the data presented in figure 5 and table 1 shows that \( \Lambda_{ei} \) with \( q_m = q_mZ \) turns out to be closer to the quantum-mechanical calculations of \( \Lambda \).

The scattering of an electron by noble gas atoms is well described by the Buckingham potential [40, 41]

\[
V(r) = -\frac{\alpha_{Xe}\epsilon^2}{2\left(r^2 + r_B^2\right)},
\]

(32)
where \( \alpha_{Xe} \) is the dipole polarizability of the target atom and \( r_B \) is a phenomenological cutoff parameter. The potential (32) describes the long-range polarization potential of the electron-atom interaction, and we took into account the short-range correlation and exchange effects by specifying the initial condition for the phase shift: \( \delta_{\ell=0}(r = 0) = k\Lambda_0 \). In calculations we set \( \Lambda_0 = -5.7a_0 \) [42]. Figure 6 compares the cross sections calculated by the PSM at different
Figure 5. The “Coulomb” logarithm (31) at $z_i = 1$, $T = 300$ K and various electron number densities: $n_e = n_i = 10^{14}$ (1), $10^{16}$ (2), $10^{18}$ (3), and $10^{20}$ cm$^{-3}$ (4), the Coulomb logarithm (13) at $q_m = q_mZ$ (5) and at $q_m = q_mS$ (6).

Figure 6. Transport cross sections for e-Xe scattering: the experimental data [43] (◦), PSM data with the Buckingham potential at $r_B = 1.168a_0$ and $\Lambda_0 = -5.7a_0$ (——), $r_B = 1.18a_0$ and $\Lambda_0 = -5.7a_0$ (- - - -), $r_B = 1.2a_0$, $\Lambda_0 = -5.7a_0$ (— —), and $r_B = 1.2a_0$ and $\Lambda_0 = 0$ (— — —).

Table 1. The Coulomb logarithm, $\Lambda_{ei}$, for $S_i = 1$ according to (13) at $T = 300$ K, $z_i = 1$ for two options of the boundary value of the wavenumber.

| $n_e$ (cm$^{-3}$) | $q_m = q_mZ$ (19) | $q_m = q_mS$ (18) |
|------------------|--------------------|--------------------|
| $10^{14}$        | 0.96               | 0.45               |
| $10^{16}$        | 0.28               | $5.1 \times 10^{-4}$ |
| $10^{18}$        | $4.2 \times 10^{-2}$ | $5.5 \times 10^{-8}$ |
| $10^{20}$        | $1.3 \times 10^{-2}$ | $3.0 \times 10^{-9}$ |

values of $r_B$ and $\Lambda_0$ with experimental data [43]. Good agreement between the calculation and experiment is seen, which makes it possible to use (32) at calculating the scattering cross section of electrons in dense noble gases, taking into account the correlation of atoms.

5. Conclusions

The method for calculating the tensors of electrical conductivity, Seebeck coefficient and thermal conductivity of a nonideal plasma in a magnetic field is described. It is shown that there is a density range in which all components of the Seebeck tensor in xenon change their sign due to the Ramsauer minimum effect in the cross section for electron scattering by neutral atoms in the range of comparable values of the cyclotron and transport frequencies of electrons.

Various methods for determining the Coulomb logarithm in the kinetic theory are considered. The correlation of ions is taken into account on the basis of the OZ integral equation in the HNC approximation. It is found that the effect of ion correlation in a nondegenerate plasma is weak, while in a degenerate plasma this effect must be taken into account. Calculations of the transport cross sections for electron scattering by hydrogen ions and xenon atoms by the phase shift method have been performed.
Acknowledgments
The work is supported by the Russian Science Foundation (grant No. 16-12-10424).

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