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Correlation and quantum effects on collisional absorption in dense laser-produced plasmas

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Abstract

Starting from quantum kinetic theory, collisional absorption of laser radiation is investigated for dense plasmas. The electrical current and the energy balance equations are formulated within the framework of nonequilibrium Green's function methods. Quantum statistical expressions are derived for plasmas in which the coupling between electrons and ions is weak due to the influence of the strong high-frequency laser field, however, the electron and the ion components may be strongly coupled within their respective subsystems. Consequently, the expressions for, e.g., the electrical current and the cycle-averaged energy absorption rate contain the dynamical structure factors and the dielectric function of the strongly correlated electron and ion subsystems. The expressions are valid for arbitrary field strength assuming the nonrelativistic case.

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1. Introduction

Due to the impressive progress in laser technology, which makes femtosecond laser pulses of very high intensity available in laboratory experiments [1], the laser–matter interaction has become a field of current interest. If the solid target is irradiated by such a laser pulse, dense plasmas can be created relevant for astrophysics and inertial confinement fusion. At high intensities the quiver velocity can be larger compared to the thermal velocity and interesting nonlinear effects have to be expected.

An important mechanism of energy deposition is inverse bremsstrahlung (IB), i.e. laser light absorption via collisional processes between the plasma particles. In strongly ionized plasmas, this absorption process is essentially governed by the electron–ion interaction usually described in terms of the electron–ion collision frequency.

In several papers, various approaches were used to calculate the electron–ion collision frequency and the dynamic conductivity, respectively, for classical plasmas under different conditions [2–5]. They cannot be applied to situations of laser–plasma interaction where quantum effects become important. Quantum effects in dense plasmas can be expected (i) if the Landau length $l = e^2/k_B T_e$ is comparable to the thermal wave length $\lambda = (2\pi\hbar^2/m_e k_B T_e)^{1/2}$, i.e. $l/\lambda \leqslant 1$, (ii) for $\hbar\omega/k_B T_e > 1$ with ω being the laser frequency, and (iii) if the electrons with number density n_e have to be described by Fermi statistics in degenerate plasmas, i.e. $n_e\lambda^3 > 1$. Quantum mechanical treatments were given by several authors [6, 7]. However rigorous quantum kinetic approaches, to the IB absorption in dense plasmas were missing until recently. Kremp *et al* [8] derived a quantum kinetic equation for dense plasmas in strong laser fields using nonequilibrium Green's function techniques. In this approach, the different interaction processes can be taken into account by appropriate approximations of the generalized field-dependent scattering rates including nonlinear field effects such as multiphoton processes and higher harmonics generation. Subsequently, quantum statistical expressions for the electron–ion collision frequency were derived, and time-dependent phenomena were studied by numerical solution of this equation [9, 10]. Quantum expressions for the collision term and the electron–ion collision frequency including dynamic screening were first given in [11–14]. The main focus of the present paper is on generalizations of the approach to study effects of strong electron–electron and ion–ion correlations on the collisional absorption rate [15, 16].

2. Kinetic equation for plasmas in electromagnetic fields

We consider a plasma under the influence of intense laser radiation. The plasma is assumed to be fully ionized consisting of electrons with number densities n_e and ions of charge $e_i = Ze$ with number density n_i. Equilibrium and nonequilibrium properties of strongly correlated plasmas are successfully described using the method of real-time Green's functions. In this framework, the nonequilibrium plasma state is given by the two-time correlation functions that are averages over creation and annihilation operators $\psi^\dagger$ and ψ,

$$g_a^>(1,1') = \frac{1}{\mathrm{i}\hbar}\langle\psi_a(1)\psi_a^\dagger(1')\rangle \qquad g_a^<(1,1') = -\frac{1}{\mathrm{i}\hbar}\langle\psi_a^\dagger(1')\psi_a(1)\rangle \tag{1}$$

where $1 \equiv \left(\mathbf{r}_1, t_1, s_1^3\right)$, and a labels the particle species. For equal times $t_1 = t_1'$, the density matrix or Wigner function, respectively, follows from the correlation function $g_a^<$. The time evolution of $g_a^{\gtrless}$, which contain the complete dynamical and statistical information, is determined by the Kadanoff–Baym equations [8]

$$\begin{aligned}\left[\mathrm{i}\hbar\frac{\partial}{\partial t_1} - \frac{1}{2m_a}\left(\frac{\hbar}{\mathrm{i}}\nabla_1 - \frac{e_a}{c}\mathbf{A}(1)\right)^2 - e_a\phi(1)\right] g_a^{\gtrless}(1,1') \\ = \sum_b \int \mathrm{d}2 V_{ab}(1-2) g_{ab}(12,1'2^+) \\ = \int \mathrm{d}\bar{1}\left[\Sigma_a^R(1,\bar{1})g_a^{\gtrless}(\bar{1},1') + \Sigma_a^{\gtrless}(1,\bar{1})g_a^A(\bar{1},1')\right]\end{aligned} \tag{2}$$

with g_{ab} being the respective two-particle Green's function and $V_{ab}(1-2) = V_{ab}(\mathbf{r}_1 - \mathbf{r}_2)\delta(t_1 - t_2)$ being the Coulomb potential. In the last line, introduction of the self-energy functions $\Sigma_a^{\gtrless}$ and Σ_a^R decouples the hierarchy of equations.

We consider in the following a spatially homogeneous electric field using vector potential gauge (**A** and ϕ denote the vector and scalar potentials, respectively),

$$\mathbf{A}(t) = -\int_{-\infty}^{t} \mathrm{d}\bar{t}\, \mathbf{E}(\bar{t}) \qquad \phi = 0. \tag{3}$$

It is advantageous to have a gauge invariant description which is achieved by the following transform [17]:

$$\tilde{g}_a(\mathbf{k}; t, t') = \int \mathrm{d}^3 r \exp\left(-\frac{\mathrm{i}}{\hbar}\mathbf{r}\cdot\left(\mathbf{k} + e_a \int_{t'}^{t} \mathrm{d}\bar{t} \frac{\mathbf{A}(\bar{t})}{t - t'}\right)\right) g_a(\mathbf{r}; t, t'). \tag{4}$$

The Wigner distribution function is given by $\tilde{f}_a(\mathbf{k}, t) = -\mathrm{i}\hbar \tilde{g}_a^<(\mathbf{k}; t, t)$.

The kinetic equation for this gauge invariant Wigner distribution function follows from the time diagonal Kadanoff–Baym equation and it reads [8]

$$\left\{\frac{\partial}{\partial t} + e_a \mathbf{E}(t)\cdot\nabla_{k_a}\right\} \tilde{f}_a(\mathbf{k}_a, t) = -2\,\mathrm{Re}\int_{t_0}^{t} \mathrm{d}\bar{t}\left\{\tilde{\Sigma}_a^> \tilde{g}_a^< - \tilde{\Sigma}_a^< \tilde{g}_a^>\right\} = \sum_b I_{ab}(\mathbf{k}_a, t) \tag{5}$$

where a and b label the particle species. The collision integral on the right-hand side (rhs) can be expressed in terms of the self-energy functions $\tilde{\Sigma}_a^{\gtrless}\left(\mathbf{k}_a + \mathbf{K}_a^A(t, \bar{t}); t, \bar{t}\right)$ and the two-time correlation functions $\tilde{g}_a^{\gtrless}\left(\mathbf{k}_a + \mathbf{K}_a^A(t, \bar{t}); \bar{t}, t\right)$, where $\mathbf{K}_a^A(t, \bar{t}) = e_a \int_{\bar{t}}^{t} \mathrm{d}t' [\mathbf{A}(t) - \mathbf{A}(t')]/(t - \bar{t})$ is the field-induced momentum shift.

The kinetic equation (5) is still very general. Two steps are necessary to find a closed form with an explicit expression for the collision term: (i) the self-energy functions or, equivalently, the two-particle Green's function, have to be specified in a certain approximation, and (ii) the two-time correlation functions $g_a^{\gtrless}$ have to be expressed in terms of the Wigner distribution function. The latter task is known as the reconstruction problem [18].

Powerful schemes are available to determine appropriate approximations for the self-energy function taking into account nonlinear field dependence as well as many-body and quantum effects relevant for high-density plasmas. In previous papers [11, 13] we have used the so-called V^s-approximation. It reads for the gauge invariant Fourier transform

$$\tilde{\Sigma}_a^{\gtrless}(\mathbf{k}; t, t') = \mathrm{i}\hbar \int \frac{\mathrm{d}^3 q}{(2\pi\hbar)^3} \tilde{g}_a^{\gtrless}(\mathbf{k} - \mathbf{q}; t, t') V_{aa}^{s\,\gtrless}(\mathbf{q}; t, t') \tag{6}$$

and after inserting this expression into equation (5) we have the starting point to calculate the properties of interest. The details of the calculation are shown in [13]. The V^s-approximation corresponds to a dynamically screened Born approximation and is therefore applicable to weakly coupled laser plasmas. In the next section, we will give a generalization to the case of strong coupling.

3. Collisional absorption for strongly correlated laser plasmas

As we are interested here in the collisional absorption by the dense plasma, it is obvious to start from the balance equation for the energy and the electrical current resulting from the kinetic equation (5). The energy balance reads

$$\frac{\mathrm{d}W^{\mathrm{kin}}}{\mathrm{d}t} - \mathbf{j}\cdot\mathbf{E} = \sum_{a,b} \int \frac{\mathrm{d}^3 k_a}{(2\pi\hbar)^3} \frac{k_a^2}{2m_a} I_{ab}(\mathbf{k}_a). \tag{7}$$

One can show that the rhs of equation (7) with a non-Markovian collision integral gives just the contribution of the mean value of the potential energy. Thus the energy balance (7) is given by

$$\frac{\mathrm{d}W^{\mathrm{kin}}}{\mathrm{d}t} + \frac{\mathrm{d}W^{\mathrm{pot}}}{\mathrm{d}t} = \mathbf{j}\cdot\mathbf{E} \tag{8}$$

where the rhs is in turn the energy loss of the electromagnetic field due to Poynting's theorem.

For the calculation of the collisional absorption, we start from the general balance equation for the current density following from equation (5) in which the collision term is not expressed in terms of the self-energy but in terms of $L_{ab} = \mathrm{i}\hbar(g_{ab} - g_a g_b)$. We find

$$\frac{\mathrm{d}}{\mathrm{d}t}\mathbf{j}_a(t) - n_a\frac{e_a^2}{m_a}\mathbf{E}(t) = \sum_{b\neq a}\int\frac{\mathrm{d}^3q}{(2\pi\hbar)^3}\frac{e_a\mathbf{q}}{m_a}V_{ab}(q)L_{ab}^{<}(\mathbf{q};t,t) \tag{9}$$

where $\mathrm{i}\hbar L_{ab}^{<}(t,t') = \langle\delta\rho_b(t')\delta\rho_a(t)\rangle$ denotes the correlation function of the density fluctuations. For this function an appropriate approximation has to be found.

The two-time correlation function L_{ab} of the density fluctuations is determined by the following general equations of motion (on the so-called Keldysh contour [19, 20], for brevity all arguments are suppressed):

$$L_{ab} = \Pi_{ab} + \sum_{c,d}\Pi_{ac}V_{cd}L_{db}. \tag{10}$$

In a plasma in a strong laser field, the coupling between species with different charges can be considered to be weak, whereas the coupling between particles with equal charges is not affected. Then the so-called polarization functions Π_{ab} can be adopted to be diagonal, $\Pi_{ab} = \delta_{ab}\Pi_a$, and an approximation in lowest order of V_{ei} is appropriate. We find

$$L_{ei}^{\gtrless}(\mathbf{q};t,t') = \int\mathrm{d}\bar{t}\big[\mathcal{L}_{ee}^{\gtrless}(\mathbf{q};t,\bar{t})V_{ei}(q)\mathcal{L}_{ii}^{A}(\mathbf{q};\bar{t},t') + \mathcal{L}_{ee}^{R}(\mathbf{q};t,\bar{t})V_{ei}(q)\mathcal{L}_{ii}^{\gtrless}(\mathbf{q};\bar{t},t')\big]. \tag{11}$$

Here the two auxiliary functions $\mathcal{L}_{ee}$ and $\mathcal{L}_{ii}$ are defined by

$$\mathcal{L}_{ee} = \Pi_e + \Pi_e V_{ee}\mathcal{L}_{ee} \qquad \mathcal{L}_{ii} = \Pi_i + \Pi_i V_{ii}\mathcal{L}_{ii}. \tag{12}$$

Consequently, the functions $\mathcal{L}_{aa}^{R/A}$ and $\mathcal{L}_{aa}^{\gtrless}$ are density response functions and correlation functions of density fluctuations, respectively, of the two subsystems.

For the electron current it follows that

$$\frac{\mathrm{d}}{\mathrm{d}t}\mathbf{j}_e(t) - n_e\frac{e_e^2}{m_e}\mathbf{E}(t) = \mathrm{Re}\int\frac{\mathrm{d}^3q}{(2\pi\hbar)^3}\frac{e_e\mathbf{q}}{m_e\mathrm{i}\hbar}V_{ei}(q)2\pi\int_{t_0}^{t}\mathrm{d}\bar{t}\big[\mathcal{S}_{ee}(\mathbf{q};t,\bar{t})V_{ei}(q)\mathcal{L}_{ii}^{A}(\mathbf{q};\bar{t},t) + \mathcal{L}_{ee}^{R}(\mathbf{q};t,\bar{t})V_{ei}(q)\mathcal{S}_{ii}(\mathbf{q};\bar{t},t)\big] \tag{13}$$

where we introduced the dynamical structure factor

$$2\pi\mathcal{S}_{aa}(\mathbf{q};t,\bar{t}) = \frac{\mathrm{i}\hbar}{2}\big[\mathcal{L}_{aa}^{>}(\mathbf{q};t,\bar{t}) + \mathcal{L}_{aa}^{<}(\mathbf{q};t,\bar{t})\big]. \tag{14}$$

The equation for the ion's current is similar.

The field dependence can be made explicit. One has to take into account that the functions $\mathcal{L}_{ee}$, e.g., are functionals of the electron correlation functions $g_e^{\gtrless}$ alone. Then one finds [15]

$$\mathcal{L}_{aa}\big(\mathbf{q};tt'\big|g_a^{\gtrless}\big) = \exp\left\{-\frac{\mathrm{i}}{\hbar}\mathbf{q}\cdot\frac{e_a}{m_a}\int_{t'}^{t}\mathrm{d}\bar{t}\mathbf{A}(\bar{t})\right\}\mathcal{L}_{aa}\big(\mathbf{q};t-t'\big|G_a^{\gtrless}\big) \tag{15}$$

where the function $\mathcal{L}_{aa}\big(\mathbf{q};t-t'\big|G_a^{\gtrless}\big)$ is a functional of field-free functions $G_a^{\gtrless}$ that are assumed to be local equilibrium functions. The Fourier transforms of $G_a^{\gtrless}$ are given by

$$-\mathrm{i}\hbar G_a^{<}(\mathbf{p};\omega) = A_a(\mathbf{p};\omega)f_a(\omega,T_a) \qquad \mathrm{i}\hbar G_a^{>}(\mathbf{p};\omega) = A_a(\mathbf{p};\omega)\,[1-f_a(\omega,T_a)] \tag{16}$$

with $f_a(\omega,T_a) \equiv 1/[\exp\{(\hbar\omega-\mu_a)/(k_BT_a)\}+1]$, and A_a being the spectral function of the strongly correlated subsystem.

The electric field occurs in equation (15) in an exponential factor and thus causes nonlinear effects such as multiphoton absorption and the occurrence of higher harmonics in the current. In the following, a harmonic electric field $\mathbf{E} = \mathbf{E}_0 \cos \omega t$ will be considered. The exponential pre-factor in (15) can be expanded into a Fourier series. The current balance is given then by

$$\frac{\mathrm{d}}{\mathrm{d}t}\mathbf{j}_e(t) - n_e \frac{e_e^2}{m_e}\mathbf{E}(t) = \mathrm{Re}\int \frac{\mathrm{d}^3 q}{(2\pi\hbar)^3}\frac{2\pi e_e \mathbf{q}}{m_e \hbar} V_{ei}^2(q) \sum_m \sum_n (-\mathrm{i})^{m+1} J_n\left(\frac{\mathbf{q}\cdot\mathbf{w}_{ei}^0}{\hbar\omega}\right) \times J_{n-m}\left(\frac{\mathbf{q}\cdot\mathbf{w}_{ei}^0}{\hbar\omega}\right) \mathrm{e}^{\mathrm{i}m\omega t} \int_{-\infty}^{\infty} \frac{\mathrm{d}\bar{\omega}}{2\pi}\big[\mathcal{S}_{ee}(\mathbf{q};\bar{\omega}-n\omega)\mathcal{L}_{ii}^A(\mathbf{q};\bar{\omega}) + \mathcal{L}_{ee}^R(\mathbf{q};\bar{\omega}-n\omega)\mathcal{S}_{ii}(\mathbf{q};\bar{\omega})\big] \tag{17}$$

where J_l is the Bessel function of lth order and with $\mathbf{w}_{ei}^0 = (e_e/m_e - e_i/m_i)\mathbf{E}_0/\omega$. In the above equation, we have suppressed (and will do in the following) the functional dependence of $\mathcal{S}_{aa}$ and $\mathcal{L}_{aa}$ on $G_a^{\gtrless}$.

The ion dynamic structure factor $\mathcal{S}_{ii}$ and the response function $\mathcal{L}_{ii}$ are localized in the low-frequency region, i.e., for a high-frequency electric field, $\bar{\omega}$ can be neglected in comparison with $n\omega$. In this case, the first term in the brackets vanishes because $\int \mathrm{d}\bar{\omega}\mathcal{L}_{ii}^A(\mathbf{q};\bar{\omega}) = 0$, and for the current it follows that

$$\mathbf{j}_e(t) - \int_{-\infty}^{t} \mathrm{d}\bar{t}\,\frac{n_e e_e^2}{m_e}\mathbf{E}(\bar{t}) = \mathrm{Re}\int \frac{\mathrm{d}^3 q}{(2\pi\hbar)^3}\sum_m \sum_n \frac{e_e}{m_e}\frac{\mathbf{q}}{m\hbar\omega} V_{ei}^2(q)(-\mathrm{i})^{m+2}\,\mathrm{e}^{\mathrm{i}m\omega t} \times J_n\left(\frac{\mathbf{q}\cdot\mathbf{w}_{ei}^0}{\hbar\omega}\right) J_{n-m}\left(\frac{\mathbf{q}\cdot\mathbf{w}_{ei}^0}{\hbar\omega}\right) \mathcal{L}_{ee}^R(\mathbf{q};-n\omega) n_i \mathcal{S}_{ii}(\mathbf{q}) \tag{18}$$

with the static structure factor $\mathcal{S}_{ii}(\mathbf{q})$ [21, 24] defined by

$$\mathcal{S}_{ii}(\mathbf{q}) \equiv \frac{1}{n_i}\int \mathrm{d}\bar{\omega}\mathcal{S}_{ii}(\mathbf{q},\bar{\omega}) = 1 + n_i \int \mathrm{d}^3 r [g_{ii}(\mathbf{r}) - 1]\,\mathrm{e}^{-\frac{\mathrm{i}}{\hbar}\mathbf{q}\cdot\mathbf{r}} \tag{19}$$

where $g_{ii}(r)$ is the pair correlation function.

The Fourier coefficients of the current can be identified easily from equation (18). Only the odd harmonics are allowed due to the symmetry of the interaction, cf [13].

Equation (18) is a generalization of the theory developed in [13]. Approximating $\mathcal{S}_{ii}(\mathbf{q}) \approx 1$ and using $\mathcal{L}_{ee}^R$ in random phase approximation (RPA), one gets the results of section IV in that former paper. Now there is included the static structure factor of the ion component. Further the function $\mathcal{L}_{ee}^R$ is the exact density response function of the electron subsystem, i.e. the electron–electron interaction can be included on a very high level. Appropriate approximations can be expressed via local field corrections (LFC), see e.g. [22, 23],

$$\mathcal{L}_{ee}^R(\mathbf{q},\omega) = \frac{\chi_e^0(\mathbf{q},\omega)}{1 - V_{ee}(q)G(q)\chi_e^0(\mathbf{q},\omega)} \tag{20}$$

with χ_e^0 being the usual free-electron Lindhard polarizability.

4. Collisional absorption rate

An important quantity is the cycle-averaged dissipation of energy $\langle \mathbf{j}\cdot\mathbf{E}\rangle$. Often also the electron–ion collision frequency ν_{ei} is discussed which is defined for the high-frequency case by (ω_p—plasma frequency)

$$\nu_{ei} = \frac{\omega^2}{\omega_\mathrm{p}^2}\frac{\langle \mathbf{j}\cdot\mathbf{E}\rangle}{\langle \epsilon_0 \mathbf{E}^2\rangle}. \tag{21}$$

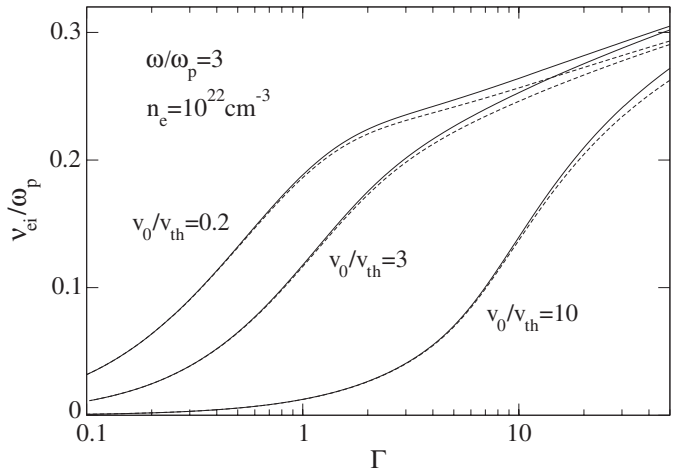

Figure 1. Electron–ion collision frequency as a function of the coupling parameter Γ for different values of the quiver velocity. LFC in accordance with Ichimaru and Utsumi [23] (solid), without LFC (dashed). S_{ii} is calculated in HNC approximation. The quiver velocity is defined as $v_0 = eE_0/m_e\omega$, $v_{\rm th} = (k_B T_e/m_e)^{1/2}$.

From equation (18) it follows that

$$\langle \mathbf{j}\cdot\mathbf{E}\rangle = \int \frac{\mathrm{d}^3q}{(2\pi\hbar)^3}\frac{\epsilon_0 q^2}{\hbar^2 e_e^2}V_{ei}^2(q)n_i\mathcal{S}_{ii}(\mathbf{q})\sum_{n=-\infty}^{+\infty} n\omega J_n^2\left(\frac{\mathbf{q}\cdot\mathbf{w}_{ei}^0}{\hbar\omega}\right)\mathrm{Im}\,\frac{1}{\varepsilon_e^R(\mathbf{q};-n\omega)}. \tag{22}$$

with $\varepsilon_e^{-1}(\mathbf{q};\omega) = 1 + \left(\hbar^2 e_e^2\right)\!/(\epsilon_0 q^2)\mathcal{L}_{ee}^R(\mathbf{q};\omega)$.

Using the dielectric function in RPA and neglecting the ion structure factor, one gets the results we derived in [13] and which were also reproduced by Kull and Plagne [14] in the framework of a quantum Vlasov approach. Those results have a similar form as that of the nonlinear Dawson–Oberman model [3]. We want to stress, however, the fact that in our result the dielectric function is given by the quantum Lindhard form, whereas the dielectric theory of Decker *et al* leads to the classical Vlasov dielectric function. The quantum effects have been discussed in [13] extensively.

Expression (22) generalizes those results to dense quantum plasmas including strong correlation effects. In the following, the influence of these correlation effects will be investigated. Figure 1 shows the influence of the LFC (the structure factor is calculated in hypernetted chain (HNC) approximation). The collision frequency is given as a function of the coupling parameter $\Gamma = (e^2/4\pi\varepsilon_0)/dk_BT$ with $d = (4\pi n_i/3)^{-1/3}$. For weak and moderate electric fields, $v_0/v_{\rm th} = 0.2$ and 3, there occur deviations in the region $\Gamma > 1$ which increase with increasing coupling. For rather strong fields, the LFC has no influence up to a coupling of about $\Gamma = 10$. Furthermore, one can see that for strong coupling the influence of the field strength becomes smaller.

Let us consider now the influence of the ion–ion correlations. Effects of such correlations via static structure factors were discussed already by Dawson and Oberman [2] starting from the linearized Vlasov equation and recently by Hazak *et al* [16] using the quantum BBGKY hierarchy. In figure 2 the collision frequency is shown using different approximations for the static ion–ion structure factor S_{ii} (here the LFC is calculated as in [23]). Inclusion of the structure factor decreases the collision frequency for small and moderate coupling up to a value of $\Gamma \approx 5$. For high values of the coupling parameter there is a strong increase in

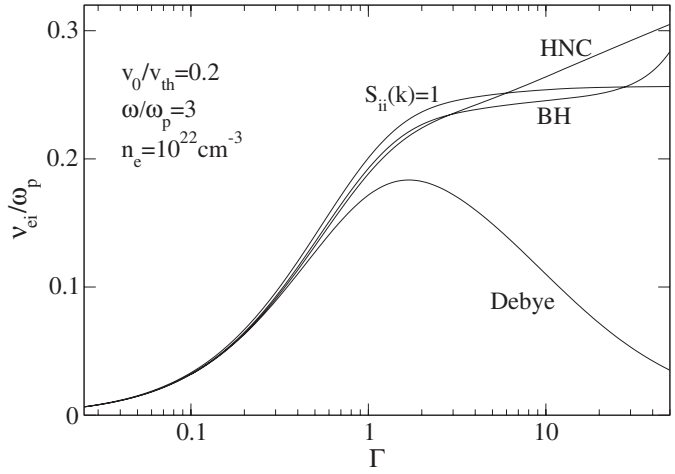

Figure 2. Electron–ion collision frequency as a function of the coupling parameter Γ. Ion structure factor in different approaches: HNC, Debye, BH (Baus–Hansen formula).

the collision frequency. This increase is even stronger in the HNC calculation than in the semi-analytical formula of Baus and Hansen [24]. As expected, the Debye approximation can be applied only for weak coupling.

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