Effect of correlation on conductivity and relaxation time

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The frequency-dependent conductivity and relaxation time of a correlated two-component plasma is calculated. We take into account the short-range correlation effect of the electron-electron, hole-hole, and electron-hole interactions by making use of the effective two-body interaction which is related to the pair correlation function. Our result for the conductivity is given in terms of the density fluctuation of the plasma. However, the present result depends on the dielectric response beyond the random-phase approximation. We take into account the effects of short-range correlations on the screening. Moreover, we consider the effect of the electron-hole short-range correlation on the electron-hole scattering-matrix element. Here we take into account the electron-hole attractive interaction during the scattering process, beyond the Born approximation, as it is commonly used.

I. INTRODUCTION

The problem of the absorption of electromagnetic waves in plasmas was studied by many authors in the 1960s.[1–5] It has become since then a well-understood problem. Here one solves the response of the electron-ion (-hole) system taking into account correctly the effect of the self-consistent field generated via the fluctuations of the charges while considering the electron-ion (-hole) correlation effects only within the Born approximation. This leads to a correct result of the absorption in the limit when the plasma parameter r_s approaches zero. This plasma parameter r_s is, in principle, given by the ratio of the average potential energy of, say the electron to its average kinetic energy. This theory met with considerable success in comparing it to experiments in wide variety of problems, such as absorption in semiconductors, metals, and nondegenerate or classical plasmas.

Similarly, much effort was directed in understanding the response of plasmas to longitudinal electric fields. To lowest order in the plasma parameter, the dielectric response depends solely on the self-consistent fields. The calculations of the dielectric function are given by the well-known random-phase approximation (RPA). However, RPA gives the correct result for the dielectric response for finite values of r_s, only when small wave numbers (i.e., large distances) are considered. For large wave numbers (i.e., short-range phenomena), one finds the RPA results unacceptable. Here one must take into account the short-range effects of the charged particles due to exchange and correlations which deviate remarkably from the RPA calculations.

The problem of strongly correlated electron plasmas has been considered by many authors.[6–16] They were mainly interested in the effect of short-range correlation on the dielectric response function and the pair correlation function. In this paper we will focus our attention on the effects of short-range correlations on the conductivity and the relaxation time of plasmas. The purpose of this paper is to develop an overall calculation scheme which adequately treats the correlation in dynamical conductivity. We shall use an equation of motion method to derive the general expression for the conductivity. In our equation of motion for the density matrix, we approximate the three-particle correlation functions by the products of one- and two-particle correlation functions, respectively. We introduce an effective interaction term instead of the Coulomb matrix element into our equations. The short-range correlation effect will be included in this effective interaction in a way which was previously suggested in Ref. 7.

II. GENERAL FORMALISM OF TOTAL CURRENT

In the following we are concerned with a many-body system described by the Hamiltonian

$$H_0=\sum_{\mathbf{k},s}E_{\mathbf{k},s}a^{\dagger}_{\mathbf{k},s}a_{\mathbf{k},s}+\frac{1}{2}\sum_{\mathbf{k},\mathbf{k}',\mathbf{q}}\sum_{s,s'}a^{\dagger}_{\mathbf{k}+\mathbf{q},s}a^{\dagger}_{\mathbf{k}'-\mathbf{q},s'}V_q^{ss'}a_{\mathbf{k}',s'}a_{\mathbf{k},s}\ , \tag{2.1}$$

in which $a^{\dagger}_{\mathbf{k},s}$ and $a_{\mathbf{k},s}$ are the usual creation and annihilation operators for state of wave vector $\mathbf{k}$ and component s, that satisfy the anticommutation relation

$$[a^{\dagger}_{\mathbf{k},s},a_{\mathbf{k}',s'}]=\delta_{\mathbf{k},\mathbf{k}'}\delta_{s,s'}\ . \tag{2.2}$$

In Eq. (2.2) $E_{\mathbf{k},s}=k^2/2m_s$. For notational convenience, we use the units in which Planck's constant $\hbar$, speed of light c, and Boltzmann's constant k_B are equal to unity through our paper. Now let us imagine that the system described by Eq. (2.1) is perturbed by an application of a weak uniform time-varying field, then

$$E_{\mathbf{k},s}\rightarrow\frac{(\mathbf{k}-e\,\mathbf{A})^2}{2m_s}\ , \tag{2.3}$$

where $\mathbf{A}$ is the vector potential of the applied field. Now the Hamiltonian of the system becomes

$$H = H_0 + H_1 \ , \tag{2.4}$$

where H_0 is still given by Eq. (2.1) and H_1 is a small perturbation

$$H_1 = \sum_{\mathbf{k},s} \mathbf{k} \cdot \mathbf{A} \frac{e_s}{m_s} a^\dagger_{\mathbf{k},s} a_{\mathbf{k},s} e^{-i\omega t} \ . \tag{2.5}$$

In writing Eq. (2.5), we have omitted the quadratic terms in **A**. For the system we are considering, the current operator is

$$\mathbf{J} = \frac{\delta H}{\delta \mathbf{A}} = \sum_{\mathbf{k},s} \frac{e_s}{m_s} (\mathbf{k} + e\,\mathbf{A}) a^\dagger_{\mathbf{k},s} a_{\mathbf{k},s} = \mathbf{J}_0 + \mathbf{J}_1 \ . \tag{2.6}$$

J_0 and J_1 are expressed as the following:

$$\mathbf{J}_0 = \sigma_0 \mathbf{E}, \quad \sigma_0 = \frac{ie^2}{\omega} \sum_s \frac{n_s}{m_s} \tag{2.7}$$

where n_s is the number density of s species of the system,

$$\mathbf{J}_1 = \sum_{\mathbf{k},s} \frac{e_s}{m_s} \mathbf{k} \langle a^\dagger_{\mathbf{k},s} a_{\mathbf{k},s} \rangle \ . \tag{2.8}$$

We see from Eq. (2.8) that the important correlation correction to the conductivity is contained in $\mathbf{J}_1$. The expectation value in Eq. (2.8) is taken with respect to the perturbed density matrix. $\mathbf{J}_1$ will have a different frequency dependence from that of $\mathbf{J}_0$ and is complex, rather than pure imaginary. This indicates that the correlation correction will affect the absorption. We now turn to this aspect of the calculation.

III. THE EVALUATION OF THE TOTAL CURRENT

The quantity we wish to evaluate is

$$\mathbf{J}_1 = \sum_{\mathbf{k},s} \frac{e_s}{m_s} \mathbf{k} \langle a^\dagger_{\mathbf{k},s} a_{\mathbf{k},s} \rangle \ , \tag{3.1}$$

which contains the expectation value of the operator $a^\dagger_{\mathbf{k},s}$ calculated with respect to the perturbed density matrix of the many-body system. If one assumes the density matrix expanded in powers of **A**, one may write

$$\rho = \rho_0 + \rho_1 + \cdots \tag{3.2}$$

and

$$\langle a^\dagger_{\mathbf{k},s} a_{\mathbf{k},s} \rangle = \mathrm{tr}(\rho_0 a^\dagger_{\mathbf{k},s} a_{\mathbf{k},s}) + \mathrm{tr}(\rho_1 a^\dagger_{\mathbf{k},s} a_{\mathbf{k},s}) = \langle a^\dagger_{\mathbf{k},s} a_{\mathbf{k},s} \rangle_0 + \langle a^\dagger_{\mathbf{k},s} a_{\mathbf{k},s} \rangle_1 \ . \tag{3.3}$$

Only the second term in Eq. (3.3) contributes to the current, so one obtains the following expression for $\mathbf{J}_1$:

$$\mathbf{J}_1 = \sum_{\mathbf{k},s} \frac{e_s}{m_s} \mathbf{k} \langle a^\dagger_{\mathbf{k},s} a_{\mathbf{k},s} \rangle_1 \ . \tag{3.4}$$

We use the equation of motion method first presented by Suhl and Werthamer[17] and by Wolff.[18] We first find the equation of motion in the Heisenberg representation of the number operator $a^\dagger_{\mathbf{k},s}(t) a_{\mathbf{k},s}(t)$. Its time derivative is determined in the usual way, by the commutator with the Hamiltonian

$$i\frac{\partial}{\partial t} a^\dagger_{\mathbf{k},s}(t) a_{\mathbf{k},s}(t) = [a^\dagger_{\mathbf{k},s}(t) a_{\mathbf{k},s}(t), H_0 + H_1] = \sum_{\mathbf{k}',\mathbf{q},s'} V_q^{ss'} e_s e_{s'} [a^\dagger_{\mathbf{k},s}(t) a^\dagger_{\mathbf{k}'-\mathbf{q},s'}(t) a_{\mathbf{k}',s'}(t) a_{\mathbf{k}-\mathbf{q},s}(t) - a^\dagger_{\mathbf{k}+\mathbf{q},s}(t) a^\dagger_{\mathbf{k}'-\mathbf{q},s'}(t) a_{\mathbf{k}',s'}(t) a_{\mathbf{k},s}(t)] \ . \tag{3.5}$$

Now it is important to notice that the equation does not explicitly contain the perturbation **A**. Thus the occupation numbers $a^\dagger_{\mathbf{k},s}(t) a_{\mathbf{k},s}(t)$ are not directly altered by the field, but only in an indirect way through its effect on the operators $a^\dagger_{\mathbf{k},s}(t) a^\dagger_{\mathbf{k}'-\mathbf{q},s'}(t) a_{\mathbf{k}',s'}(t) a_{\mathbf{k}-\mathbf{q},s}(t)$ and $a^\dagger_{\mathbf{k}+\mathbf{q},s}(t) a^\dagger_{\mathbf{k}'-\mathbf{q},s'}(t) a_{\mathbf{k}',s'}(t) a_{\mathbf{k},s}(t)$ which describe pair correlation in the plasma. By taking the trace of Eq. (3.5) with the first-order correction ρ_1 to the density matrix and substitute it into Eq. (3.4), we may now obtain the expression for $\mathbf{J}_1$ as

$$\mathbf{J}_1 = \frac{1}{\omega} \sum_{\mathbf{k},\mathbf{k}',\mathbf{q}} \sum_{s,s'} V_q^{ss'} e_s e_{s'} \mathbf{k} \frac{e_s}{m_s} (\langle a^\dagger_{\mathbf{k},s} a^\dagger_{\mathbf{k}'-\mathbf{q},s'} a_{\mathbf{k}',s'} a_{\mathbf{k}-\mathbf{q},s} \rangle_1 - \langle a^\dagger_{\mathbf{k}+\mathbf{q},s} a^\dagger_{\mathbf{k}'-\mathbf{q},s'} a_{\mathbf{k}',s'} a_{\mathbf{k},s} \rangle_1) \ , \tag{3.6}$$

where we have used the fact that all the first-order expectation values vary as $e^{i\omega t}$ since the perturbation is harmonic in time. In the first term of Eq. (3.6), let $\mathbf{k} \to \mathbf{k} + \mathbf{q}$, we obtain

$$\mathbf{J}_1 = \frac{1}{\omega} \sum_{\mathbf{k},\mathbf{k}',\mathbf{q}} \sum_{s,s'} V_q^{ss'} e_s e_{s'} \mathbf{q} \frac{e_s}{m_s} \langle a^\dagger_{\mathbf{k}+\mathbf{q},s} a^\dagger_{\mathbf{k}'-\mathbf{q},s'} a_{\mathbf{k}',s'} a_{\mathbf{k},s} \rangle_1 \ . \tag{3.7a}$$

It is instructive to see that $\mathbf{J}_1$ is composed of electron as well as hole contribution. We write

$$\mathbf{J}_1 = \mathbf{J}_1^e + \mathbf{J}_1^h \ , \tag{3.7b}$$

with

$$\mathbf{J}_1^e = \frac{1}{\omega} \sum_{\mathbf{k},\mathbf{k}',\mathbf{q}} \sum_{s'} V_q^{es'} \frac{e^2}{m_e} e_{s'} \mathbf{q} \langle a^\dagger_{\mathbf{k}+\mathbf{q},e} a^\dagger_{\mathbf{k}'-\mathbf{q},s'} a_{\mathbf{k}',s'} a_{\mathbf{k},e} \rangle_1 \ , \tag{3.7c}$$

and

$$\mathbf{J}_1^h=\frac{1}{\omega}\sum_{\mathbf{k},\mathbf{k}',\mathbf{q}}\sum_{s'}V_q^{hs'}\frac{e^2}{m_h}e_{s'}\mathbf{q}\langle a^\dagger_{\mathbf{k}+\mathbf{q},h}a^\dagger_{\mathbf{k}'-\mathbf{q},s'}a_{\mathbf{k}',s'}a_{\mathbf{k},h}\rangle_1 , \tag{3.7d}$$

being the electron and hole current, respectively. Going back to Eq. (3.7a), we make the transformations $s\leftrightarrow s'$, $\mathbf{k}\leftrightarrow\mathbf{k}'$, and $\mathbf{q}\to-\mathbf{q}$. We then use this new form for $\mathbf{J}_1$ together with the original one, Eq. (3.7a), to obtain an equivalent expression for $\mathbf{J}_1$ which demonstrates the full symmetry between s and s', i.e.,

$$\mathbf{J}_1=\frac{1}{2\omega}\sum_{\mathbf{k},\mathbf{k}',\mathbf{q}}\sum_{s,s'}V_q^{ss'}e_se_{s'}\mathbf{q}\left[\frac{e_s}{m_s}-\frac{e_{s'}}{m_{s'}}\right]\times\langle a^\dagger_{\mathbf{k}+\mathbf{q},s}a^\dagger_{\mathbf{k}'-\mathbf{q},s'}a_{\mathbf{k}',s'}a_{\mathbf{k},s}\rangle_1 . \tag{3.8}$$

If we use

$$\langle a^\dagger_{\mathbf{k}+\mathbf{q},s}a^\dagger_{\mathbf{k}'-\mathbf{q},s'}a_{\mathbf{k}',s'}a_{\mathbf{k},s}\rangle=\langle a^\dagger_{\mathbf{k}+\mathbf{q},s}a_{\mathbf{k},s}a^\dagger_{\mathbf{k}'-\mathbf{q},s'}a_{\mathbf{k}',s'}\rangle-\langle n_s(\mathbf{k}+\mathbf{q})\rangle\delta_{s,s'}\delta_{\mathbf{k},\mathbf{k}'-\mathbf{q}} , \tag{3.9}$$

where $n_s(\mathbf{k})=a^\dagger_{\mathbf{k}s}a_{\mathbf{k}s}$, Eq. (3.8) becomes

$$\mathbf{J}_1=\frac{1}{2\omega}\sum_{\mathbf{k},\mathbf{k}',\mathbf{q}}\sum_{s,s'}V_q^{ss'}e_se_{s'}\mathbf{q}\left[\frac{e_s}{m_s}-\frac{e_{s'}}{m_{s'}}\right]\times\langle a^\dagger_{\mathbf{k}+\mathbf{q},s}a_{\mathbf{k},s}a^\dagger_{\mathbf{k}'-\mathbf{q},s'}a_{\mathbf{k}',s'}\rangle_1 . \tag{3.10}$$

Because of the δ function $\delta_{ss'}$, the second term of Eq. (3.9) does not contribute to the current. In Eq. (3.10), the summation over $\mathbf{k},\mathbf{k}'$ can be easily completed and we obtain

$$\mathbf{J}_1=\frac{1}{2\omega}\sum_{\mathbf{k},\mathbf{k}',\mathbf{q}}\sum_{s,s'}V_q^{ss'}e_se_{s'}\mathbf{q}\left[\frac{e_s}{m_s}-\frac{e_s}{m_{s'}}\right]\langle n_\mathbf{q}(s)n_{-\mathbf{q}}(s')\rangle , \tag{3.11}$$

where

$$n_\mathbf{q}(s)=\sum_\mathbf{k}a^\dagger_{\mathbf{k}+\mathbf{q},s}a_{\mathbf{k},s} . \tag{3.12}$$

We see that the induced current is uniquely determined by correlation of induced density $\langle n_\mathbf{q}(s)n_{-\mathbf{q}}(s')\rangle$, Eq. (3.10) is an exact expression for $\mathbf{J}_1$. Its form shows immediately that there is no absorption if the system only contains one species. To determine the current $\mathbf{J}_1$ we must evaluate the perturbed correlation function $\langle a^\dagger_{\mathbf{k}+\mathbf{q},s}a_{\mathbf{k},s}a^\dagger_{\mathbf{k}'-\mathbf{q},s'}a_{\mathbf{k}',s'}\rangle_1$. We consider this question in Sec. IV.

IV. CALCULATION OF PERTURBED CORRELATION FUNCTION

To determine $\mathbf{J}_1$, we need to know $\langle n_\mathbf{q}(s)n_{-\mathbf{q}}(s')\rangle$. If one writes the equation of motion for $n_\mathbf{q}(s,t)n_{-\mathbf{q}}(s',t)$ (i.e., calculates the first time derivative of density fluctuation), all terms which are proportional to the Coulomb interaction cancel out. Therefore we shall solve for the function $\langle a^\dagger_{\mathbf{k}+\mathbf{q},s}a_{\mathbf{k},s}a^\dagger_{\mathbf{k}'-\mathbf{q},s'}a_{\mathbf{k}',s'}\rangle_1$, and then sum over k and k' to obtain $\langle n_\mathbf{q}(s)n_{-\mathbf{q}}(s')\rangle$. We consider the equation of motion of

$$\langle a^\dagger_{\mathbf{k}+\mathbf{q},s}(0)a_{\mathbf{k},s}(0)a^\dagger_{\mathbf{k}'-\mathbf{q},s'}(t)a_{\mathbf{k}',s'}(t)\rangle ,$$

which is the two-particle autocorrelation function, and in the limit of $t\to0$ gives us the desired two-particle correlation function. We shall demonstrate that the perturbed autocorrelation function obeys, in the random-phase approximation, a solvable integral equation. On the contrary, the equation of motion for

$$\langle a^\dagger_{\mathbf{k}+\mathbf{q},s}(t)a_{\mathbf{k},s}(t)a^\dagger_{\mathbf{k}'-\mathbf{q},s'}(t)a_{\mathbf{k}',s'}(t)\rangle_1$$

is very involved. A straightforward computation shows that the equation of motion for

$$\langle a^\dagger_{\mathbf{k}+\mathbf{q},s}(t)a_{\mathbf{k},s}(t)a^\dagger_{\mathbf{k}'-\mathbf{q},s'}(t)a_{\mathbf{k}',s'}(t)\rangle_1$$

is given by

$$\begin{aligned}i\frac{\partial}{\partial t}[a^\dagger_{\mathbf{k}+\mathbf{q},s}(0)a_{\mathbf{k},s}(0)a^\dagger_{\mathbf{k}'-\mathbf{q},s'}(t)a_{\mathbf{k}',s'}(t)]&=a^\dagger_{\mathbf{k}+\mathbf{q},s}(0)a_{\mathbf{k},s}(0)[a^\dagger_{\mathbf{k}'-\mathbf{q},s'}(t)a_{\mathbf{k}',s'}(t),H]\\&=(E_{\mathbf{k}',s'}-E_{\mathbf{k}'-\mathbf{q},s'})[a^\dagger_{\mathbf{k}+\mathbf{q},s}(0)a_{\mathbf{k},s}(0)a^\dagger_{\mathbf{k}'-\mathbf{q},s'}(t)a_{\mathbf{k}',s'}(t)]\\&\quad-\frac{ie_{s'}}{m_{s'}}\frac{\mathbf{q}\cdot\mathbf{E}}{\omega}e^{i\omega t}[a^\dagger_{\mathbf{k}+\mathbf{q},s}(0)a_{\mathbf{k},s}(0)a^\dagger_{\mathbf{k}'-\mathbf{q},s'}(t)a_{\mathbf{k}',s'}(t)]\\&\quad+\sum_{\mathbf{k}_1,\mathbf{q}_1,s_1}V_{q_1}e_{s'}e_{s_1}[a^\dagger_{\mathbf{k}+\mathbf{q},s}(0)a_{\mathbf{k},s}(0)a^\dagger_{\mathbf{k}'-\mathbf{q},s'}(t)a^\dagger_{\mathbf{k}_1-\mathbf{q}_1,s_1}(t)a_{\mathbf{k}_1,s_1}(t)a_{\mathbf{k}'-\mathbf{q}_1,s'}(t)]\\&\quad-\sum_{\mathbf{k}_1,\mathbf{q}_1,s_1}V_{q_1}e_{s'}e_{s_1}[a^\dagger_{\mathbf{k}+\mathbf{q},s}(0)a_{\mathbf{k},s}(0)a^\dagger_{\mathbf{k}'-\mathbf{q}+\mathbf{q}_1,s'}(t)a_{\mathbf{k}_1-\mathbf{q}_1,s_1}(t)a_{\mathbf{k}_1,s_1}(t)a_{\mathbf{k}',s'}(t)] .\end{aligned} \tag{4.1}$$

This is a set of equations which, through the Coulomb interaction, couples the two- and three-particle correlation functions. Progress in solving this set of equations is usually made by replacing the three-particle correlations by a suitable product of one- and two-particle correlations. For weakly coupled plasmas, i.e., when the average potential energy of a pair of charged particles is small compared to their kinetic energy, we use the RPA approximation. Here we retain only those terms which arise from the direct Coulomb interaction—that is, these terms in Eq. (4.1) in which **q** is equal to **q**′. These introduce a factor of $1/q^2$ into the Coulomb terms of Eq. (4.1) which are highly divergent in the limit $q \to 0$. Other terms which are not singular and thus do not affect the long-range behavior arising from the Coulomb potential are omitted. Thus the basic approximation of RPA (Ref. 18) is to replace the average of the three-particle operator by the following expression:

$$\langle a^\dagger_{\mathbf{k}+\mathbf{q},s}(0)a_{\mathbf{k},s}(0)a^\dagger_{\mathbf{k}'-\mathbf{q},s'}(t)a^\dagger_{\mathbf{k}_1-\mathbf{q}_1,s_1}(t)a_{\mathbf{k}_1,s_1}(t)a_{\mathbf{k}'-\mathbf{q}_1,s'}(t)\rangle_1 \approx \langle a^\dagger_{\mathbf{k}+\mathbf{q},s}(0)a_{\mathbf{k},s}(0)a^\dagger_{\mathbf{k}_1-\mathbf{q}_1,s_1}(t)a_{\mathbf{k}_1,s_1}(t)\rangle_1 \langle a^\dagger_{\mathbf{k}'-\mathbf{q},s'}(t)a_{\mathbf{k}'-\mathbf{q}_1,s'}(t)\rangle_0 \,. \tag{4.2}$$

We now multiply Eq. (4.1) by the density matrix, take its trace, and collect terms of first order in **A**, and make use of the approximation outlined in Eq. (4.2). We thus obtain a relatively tractable equation for the perturbation of the two-particle correlation function. Let us define

$$F^0_q(\mathbf{k},\mathbf{k}',s,s',t)=\langle a^\dagger_{\mathbf{k}+\mathbf{q},s}(0)a_{\mathbf{k},s}(0)a^\dagger_{\mathbf{k}'-\mathbf{q},s'}(t)a_{\mathbf{k}',s'}(t)\rangle_0 \,, \tag{4.3}$$

and

$$F^1_q(\mathbf{k},\mathbf{k}',s,s',t)=\langle a^\dagger_{\mathbf{k}+\mathbf{q},s}(0)a_{\mathbf{k},s}(0)a^\dagger_{\mathbf{k}'-\mathbf{q},s'}(t)a_{\mathbf{k}',s'}(t)\rangle_1 \,. \tag{4.4}$$

The equation of motion for the correlation function F^1 now reads

$$\left[i\frac{\partial}{\partial t}+E_{s'}(\mathbf{k}'-\mathbf{q})-E_{s'}(\mathbf{k}')\right]F^1_q(\mathbf{k},\mathbf{k}',s,s',t)=\frac{ie_{s'}}{m_{s'}}\frac{\mathbf{q}\cdot\mathbf{E}}{\omega}e^{i\omega t}F^0_q(\mathbf{k},\mathbf{k}',s,s',t) -[n_{s'}(\mathbf{k}'-\mathbf{q})-n_{s'}(\mathbf{k}')]\sum_{\mathbf{k}'',s''}V_q^{s's''}e_{s''}F^1_q(\mathbf{k},\mathbf{k}'',s,s'',t) \,. \tag{4.5}$$

In this RPA approximation we have taken into account the long-range effect of the Coulomb interaction. We now proceed to improve the RPA result as given by Eq. (4.5), to include the local field correction due to short-range correlation. In order to determine how to include short-range correlation we first consider the second time derivative of the operator $n_{\mathbf{q},s}(0)n_{-\mathbf{q},s'}(t)$, we find that

$$\begin{aligned}\frac{\partial^2}{\partial t^2}n_{\mathbf{q},s}(0)n_{-\mathbf{q},s'}(t)=&-n_{\mathbf{q},s}(0)\sum_{\mathbf{k}}\left[\frac{\mathbf{k}\cdot\mathbf{q}}{m_{s'}}-\frac{q^2}{m_{s'}}\right]^2 a^\dagger_{\mathbf{k}-\mathbf{q},s'}(t)a_{\mathbf{k},s'}(t)\\ &-n_{\mathbf{q},s}(0)\sum_{s_1}V_q^{s',s_1}\frac{n_{s_1}q^2}{m_{s_1}}n_{-\mathbf{q},s'}(t)-n_{\mathbf{q},s}(0)\sum_{s_1}\sum_{q'\neq q}V_{q'}^{s',s_1}\frac{\mathbf{q}\cdot\mathbf{q}'}{m_{s_1}}n_{-\mathbf{q},s'}(t)n_{\mathbf{q}'-\mathbf{q},s_1}(t)\\ &-\frac{e^2}{c}n_{\mathbf{q},s}(0)\mathbf{q}\cdot\mathbf{A}e^{i\omega t}\sum_{\mathbf{k}}\left[\frac{\mathbf{k}\cdot\mathbf{q}}{m_{s'}}-\frac{q^2}{m_{s'}}\right]^2 a^\dagger_{\mathbf{k}-\mathbf{q},s'}(t)a_{\mathbf{k},s'}(t) \,.\end{aligned} \tag{4.6a}$$

We shall first consider Eq. (4.6a) in the absence of the external field.[7] The first term of the right-hand side (rhs) of the equation represents the single-particle recoil and the Doppler shift. The second term on the rhs is due to the long-range part of the Coulomb potential and is proportional to the plasma frequency. The third term on the rhs involves the product of two density fluctuation operators. Since $n_{q-q'}(s)$, given by $\sum_i e^{i(\mathbf{q}-\mathbf{q}')\cdot\mathbf{r}_i}$ (where $\mathbf{r}_i$ represents the electronic coordinate of the ith electrons), is a sum of complex exponential terms with differing phases for $\mathbf{q}\neq\mathbf{q}'$, and since the ensemble average of $n_{q-q'}(s)$ vanishes for $(q\neq q')$ if the system is homogeneous, we expect destructive interference to occur in this term and therefore we drop it from the equation. This is the original random-phase approximation as proposed by Bohm and Pines.[16] Therefore within the RPA and without external potential, we obtain

$$\begin{aligned}\frac{\partial^2}{\partial t^2}n_{\mathbf{q}}(s,0)n_{-\mathbf{q}}(s',t) &=-n_{\mathbf{q},s}(0)\sum_{\mathbf{k}}\left[\frac{\mathbf{k}\cdot\mathbf{q}}{m_{s'}}-\frac{q^2}{m_{s'}}\right]^2 a^\dagger_{\mathbf{k}-\mathbf{q},s'}(t)a_{\mathbf{k},s'}(t)\\ &\quad -n_{\mathbf{q},s}(0)\sum_{s_1}\frac{e_s}{e_{s_1}}\omega_p^2(s_1)n_{\mathbf{q}}(s,0)n_{-\mathbf{q}}(s_1,t) \,,\end{aligned} \tag{4.6b}$$

where $\omega_p^2(s)=4\pi e^2 n_s/m_s$ is the plasma frequency for the s component. It is obvious that for the limit of $q\to 0$, the first term on the rhs of Eq. (4.6b) vanishes. By taking the ensemble average over the density operators we are left with four homogeneous coupled equations for $\langle n_{\mathbf{q}}(s,0)n_{-\mathbf{q}}(s',t)\rangle$ with one finite frequency root at $\omega^2=\omega_p^2(e)+\omega_p^2(h)$. To include the short-range correlation, the nonlinear term in Eq. (4.6a) (i.e., the third term on the rhs) must be retained. If we use the expression for $n_{\mathbf{q},s}$ in terms of the electronic coordinates $\mathbf{r}_i$, i.e., $n_{\mathbf{q},s}=\sum_i e^{i\mathbf{q}\cdot\mathbf{r}_i}$, we may rewrite this nonlinear term as

$$-n_{\mathbf{q},s}\sum_{s_1}\sum_{q'}V_q^{s's_1}\frac{\mathbf{q}\cdot\mathbf{q}'}{m_{s'}}n_{-\mathbf{q},s'}(t)n_{\mathbf{q}'-\mathbf{q},s_1}(t)$$

$$=-n_{\mathbf{q},s}\sum_{s_1}\sum_{q'}V_q^{s',s_1}\frac{\mathbf{q}\cdot\mathbf{q}'}{m_{s'}}n_{-\mathbf{q},s'}(t)n_{\mathbf{q}'-\mathbf{q},s_1}(t)n_{\mathbf{q}-\mathbf{q}',s_1}(t)\ . \tag{4.7}$$

We now follow the approximation used in Ref. 7 and replace the operators $n_{\mathbf{q}'-\mathbf{q},s_1}(t)n_{\mathbf{q}-\mathbf{q}',s_1}(t)$ in Eq. (4.7) by their ensemble average $\langle n_{\mathbf{q}'-\mathbf{q},s_1}(t)n_{\mathbf{q}-\mathbf{q}',s_1}(t)\rangle$. Since both density operators depend on the same time t, the average value of $n_{\mathbf{q}'-\mathbf{q},s_1}(t)n_{\mathbf{q}-\mathbf{q}',s_1}(t)$ is time independent and is given by the static structure factor[19,20] $S^{ss_1}(\mathbf{q}-\mathbf{q}')$. Using this approximation we can now combine the second and the third terms of Eq. (4.6a) [using Eq. (4.7)] to give

$$-n_{\mathbf{q},s}(0)\sum_{s_1}V_q^{s',s_1}\frac{n_{s_1}q^2}{m_{s_1}}n_{-\mathbf{q},s'}(t)-n_{\mathbf{q},s}(0)\sum_{s_1}\sum_{q'(\neq q)}V_{q'}^{s',s_1}\frac{\mathbf{q}\cdot\mathbf{q}'}{m_{s_1}}n_{-\mathbf{q}',s'}(t)n_{\mathbf{q}'-\mathbf{q},s_1}(t)=-n_{\mathbf{q},s}\sum_{s_1}U_q^{s',s_1}\frac{n_{s_1}q^2}{m_{s'}}n_{-\mathbf{q},s'}(t)\ ,$$

where $U_q^{s,s'}$ is the effective interaction including short-range correlations and is given by

$$U_q^{ss'}=V_q^{ss'}\left[1+\int d\mathbf{q}'\frac{\mathbf{q}\cdot\mathbf{q}'}{q^2}[S^{ss'}(\mathbf{q}-\mathbf{q}')-1]\right]\ . \tag{4.8}$$

Thus Eq. (4.6a) is approximated by

$$\frac{\partial^2}{\partial t^2}n_{\mathbf{q}}(s,0)n_{-\mathbf{q}}(s',t)=-n_{\mathbf{q},s}(0)\sum_{\mathbf{k}}\left[\frac{\mathbf{k}\cdot\mathbf{q}}{m_{s'}}-\frac{q^2}{m_{s'}}\right]^2a^\dagger_{\mathbf{k}-\mathbf{q},s'}(t)a_{\mathbf{k},s'}(t)-n_{\mathbf{q},s}(0)\sum_{s_1}\frac{e_{s'}}{e_{s_1}}\omega_p^2(s_1)\frac{U_q^{s',s_1}}{V_q^{s',s_1}}n_{\mathbf{q}}(s,0)n_{-\mathbf{q}}(s_1,t)\ . \tag{4.9}$$

Therefore we see that the equation of motion for $n_{\mathbf{q},s}(t)n_{-\mathbf{q},s'}(t)$ has the same structure as in the RPA case except that V_q, the Coulomb interaction, has been replaced by U_q, the effective interaction. We now argue that in order to obtain short-range correlation effects in Eq. (4.5), the interaction $V_q^{ss'}$ should be similarly replaced by $U_q^{ss'}$, the effective interaction, while all other terms remain unchanged. The derivation of the effective interaction and its use for the one-particle kinetic equation, in calculations of the dielectric response, was first proposed by Singwi *et al.*[7] with considerable success. It was generalized for multicomponent plasmas in Refs. 19 and 20. Here we use the same idea for the two-particle autocorrelation function. We point out that Refs. 7 and 19 deal with the response of time-dependent longitudinal fields having wave number **q**. We, on the other hand, calculate the response of plasmas to time-dependent homogeneous field (photons). The absorption rate in Refs. 7 and 19 is due to Landau damping. We go beyond this to obtain the collisional absorption which is a higher order in the plasma parameter r_s. Physically it implies that the longitudinal field supplies both the energy and momentum to the excited plasma while the photon supplies only the energy. The momentum is obtained by the electrons, say, interacting with the density fluctuation of the holes. We now write

$$\left[i\frac{\partial}{\partial t}+E_{s'}(\mathbf{k}'-\mathbf{q})-E_{s'}(\mathbf{k}')\right]F_q^1(\mathbf{k},\mathbf{k}',s,s',t)=\frac{ie_{s'}}{m_{s'}}\frac{\mathbf{q}\cdot\mathbf{E}}{\omega}e^{i\omega t}F_q^0(\mathbf{k},\mathbf{k}',s,s',t)$$
$$-[n_{s'}(\mathbf{k}'-\mathbf{q})-n_{s'}(\mathbf{k}')]\sum_{\mathbf{k}'',s''}U_q^{s's''}e_{s''}F_q^1(\mathbf{k},\mathbf{k}'',s,s'',t)\ . \tag{4.10}$$

In order to solve F^1 in terms of F^0 we use the Fourier transformations

$$F_q^1(t)=e^{-i\omega t}T_q(t)\ , \tag{4.11}$$

$$T_q(x)=\frac{1}{2\pi}\int dt\, e^{ixt}T_q(t)\ , \tag{4.12}$$

$$F_q^0(x)=\frac{1}{2\pi}\int dt\, e^{ixt}F_q^0(t)\ , \tag{4.13}$$

we obtain

$$[\omega+x+E_{s'}(\mathbf{k}'-\mathbf{q})-E_{s'}(\mathbf{k}')]T_q(\mathbf{k},\mathbf{k}',s,s',x)=\frac{ie_{s'}}{m_{s'}}\frac{\mathbf{q}\cdot\mathbf{E}}{\omega}F_q^0(\mathbf{k},\mathbf{k}',s,s',x)$$
$$-[n_{s'}(\mathbf{k}'-\mathbf{q})-n_{s'}(\mathbf{k}')]\sum_{\mathbf{k}'',s''}U_q^{s's''}e_{s''}T_q(\mathbf{k},\mathbf{k}'',s,s'',x)\ . \tag{4.14}$$

Now we define

$$\sum_{\mathbf{k}'} T_q(\mathbf{k},\mathbf{k}',s,s',x)=\tilde{T}_q(\mathbf{k},s,s',x) , \tag{4.15}$$

and write explicitly the two coupled equations for $T_q(\mathbf{k},\mathbf{k}'',e,e,x)$ and $T_q(\mathbf{k},\mathbf{k}'',e,h,x)$,

$$\tilde{T}(e\text{-}e)=eQ[U^{e\text{-}e}e\tilde{T}(e\text{-}e)+e'U^{e\text{-}h}\tilde{T}(e\text{-}h)]+A(e\text{-}e) , \tag{4.16}$$

$$\tilde{T}(e\text{-}h)=e'Q[U^{h\text{-}e}e\tilde{T}(e\text{-}e)+e'U^{h\text{-}h}\tilde{T}(e\text{-}h)]+A(e\text{-}h) , \tag{4.17}$$

where

$$Q_s(q,\omega)=\sum_{\mathbf{k}} \frac{n_s(\mathbf{k})-n_s(\mathbf{k}-\mathbf{q})}{x+\omega-E_s(\mathbf{k})-E_s(\mathbf{k}-\mathbf{q})} . \tag{4.18}$$

For simplicity we will use $Q_q(\omega)$ for $Q_e(q,\omega)$ and $B_q(\omega)$ for $Q_h(q,\omega)$. $U^{ss'}$ is an effective interaction tensor. In Eqs. (4.16) and (4.17) the dependences on $\mathbf{k}$ and q are suppressed. e is the charge for electron while e' is the charge of the hole. The quantity $A(ss')$ in Eq. (4.16) is given as

$$A_q(\mathbf{k},s,s')=\frac{ie_{s'}}{m_{s'}}\frac{\mathbf{q}\cdot\mathbf{E}}{\omega}\sum_{\mathbf{k}'}\frac{F_q^0(\mathbf{k},\mathbf{k}',s,s',x)}{x+\omega-E_s(\mathbf{k})+E_s(\mathbf{k}'-\mathbf{q})} . \tag{4.19}$$

From Eqs. (4.16) and (4.17) we obtain the solutions for $\tilde{T}(e\text{-}e)$ and $\tilde{T}(e\text{-}i)$

$$\tilde{T}(e\text{-}e)=\frac{A(e\text{-}e)(1-e^2U^{h\text{-}h}B)+ee'QU^{e\text{-}h}A(e\text{-}h)}{\tilde{\epsilon}_q(\omega+x)} , \tag{4.20}$$

and

$$\tilde{T}(e\text{-}h)=\frac{A(e\text{-}h)(1-e^2U^{e\text{-}e}Q)+ee'BU^{h\text{-}e}A(e\text{-}e)}{\tilde{\epsilon}_q(\omega+x)} , \tag{4.21}$$

where

$$\tilde{\epsilon}_q(x)=[1-e^2U_q^{e\text{-}e}Q_q(x)][1-e^2U_q^{h\text{-}h}B_q(x)] -e^4U_q^{e\text{-}h}U_q^{h\text{-}e}Q_q(x)B_q(x) . \tag{4.22}$$

We point out that this $\tilde{\epsilon}$ is not the dielectric response function. However the modes of the collective excitation are given by zeros of this $\tilde{\epsilon}$. Similarly,

$$\tilde{T}(h\text{-}h)=\frac{A(h\text{-}h)(1-e^2U^{e\text{-}e}Q)+ee'BU^{h\text{-}e}A(h\text{-}e)}{\tilde{\epsilon}_q(\omega+x)} , \tag{4.23}$$

and

$$\tilde{T}(h\text{-}e)=\frac{A(h\text{-}e)(1-e^2U^{h\text{-}h}B)+ee'QU^{e\text{-}h}A(h\text{-}h)}{\tilde{\epsilon}_q(\omega+x)} . \tag{4.24}$$

Substituting Eqs. (4.20)–(4.24) into Eq. (4.10), we obtain

$$T_q(\mathbf{k},\mathbf{k}',e\text{-}e)=\frac{ie}{m}\frac{\mathbf{q}\cdot\mathbf{E}}{\omega}\frac{F_q^0(\mathbf{k},\mathbf{k}',e\text{-}e)}{x+\omega-E_e(\mathbf{k}')+E_e(\mathbf{k}'-\mathbf{q})} +e\frac{n_e(\mathbf{k}')-n_e(\mathbf{k}'-\mathbf{q})}{x+\omega-E_e(\mathbf{k}')+E_e(\mathbf{k}'-\mathbf{q})}[eU_q^{e\text{-}e}\tilde{T}_q(\mathbf{k},e\text{-}e)+e'U_q^{e\text{-}h}\tilde{T}_q(\mathbf{k},e\text{-}h)] , \tag{4.25}$$

and

$$T_q(\mathbf{k},\mathbf{k}',e\text{-}h)=\frac{ie'}{m'}\frac{\mathbf{q}\cdot\mathbf{E}}{\omega}\frac{F_q^0(\mathbf{k},\mathbf{k}',e\text{-}e)}{x+\omega-E_h(\mathbf{k}')+E_h(\mathbf{k}'-\mathbf{q})} +e'\frac{n_h(\mathbf{k}')-n_h(\mathbf{k}'-\mathbf{q})}{x+\omega-E_h(\mathbf{k}')+E_h(\mathbf{k}'-\mathbf{q})}[eU_q^{e\text{-}e}\tilde{T}_q(\mathbf{k},e\text{-}e)+e'U_q^{h\text{-}h}\tilde{T}_q(\mathbf{k},e\text{-}h)] , \tag{4.26}$$

and similar expressions for $T_q(h\text{-}h)$ and $T_q(h\text{-}e)$ can be obtained as

$$T_q(\mathbf{k},\mathbf{k}',h\text{-}h)=\frac{ie'}{m'}\frac{\mathbf{q}\cdot\mathbf{E}}{\omega}\frac{F_q^0(\mathbf{k},\mathbf{k}',h\text{-}h)}{x+\omega-E_h(\mathbf{k}')+E_h(\mathbf{k}'-\mathbf{q})} +e'\frac{n_h(\mathbf{k}')-n_h(\mathbf{k}'-\mathbf{q})}{x+\omega-E_h(\mathbf{k}')+E_h(\mathbf{k}'-\mathbf{q})}[e'U_q^{h\text{-}h}\tilde{T}_q(\mathbf{k},h\text{-}h)+eU_q^{h\text{-}e}\tilde{T}_q(\mathbf{k},h\text{-}e)] , \tag{4.25'}$$

and

$$T_q(\mathbf{k},\mathbf{k}',h\text{-}e)=\frac{ie}{m}\frac{\mathbf{q}\cdot\mathbf{E}}{\omega}\frac{F_q^0(\mathbf{k},\mathbf{k}',h\text{-}h)}{x+\omega-E_e(\mathbf{k}')+E_e(\mathbf{k}'-\mathbf{q})}+e\frac{n_e(\mathbf{k}')-n_e(\mathbf{k}'-\mathbf{q})}{x+\omega-E_e(\mathbf{k}')+E_e(\mathbf{k}'-\mathbf{q})}[e'U_q^{e\text{-}h}\tilde{T}_q(\mathbf{k},h\text{-}h)+eU_q^{e\text{-}e}\tilde{T}_q(\mathbf{k},h\text{-}e)] . \tag{4.26'}$$

The function F_q^1 is given by Eq. (4.11). Substitute it into Eq. (3.10) and sum over s and s', we obtain the conductivity as

$$\mathbf{J}_1=\frac{1}{2\omega}\sum_{\mathbf{q}} V_q ee'\mathbf{q}\left[\frac{e}{m}-\frac{e'}{m'}\right]\sum_{\mathbf{k},\mathbf{k}',\mathbf{q}}\int_0^\infty dx[T_q(\mathbf{k},\mathbf{k}',e\text{-}h)-T_q(\mathbf{k},\mathbf{k}',h\text{-}e)] . \tag{4.27}$$

Now our task is to evaluate the zeroth-order two-particle correlation function $F^0(\mathbf{k},\mathbf{k}',s,s',x)$. For this purpose it is convenient to consider the retarded Green's function defined by

$$G_r(t)=-i\Theta(t)\langle[a^\dagger_{\mathbf{k}'-\mathbf{q},s'}(t)a_{\mathbf{k}',s'}(t),a^\dagger_{\mathbf{k}+\mathbf{q},s}(0)a_{\mathbf{k},s}(0)]\rangle . \tag{4.28}$$

The advanced Green's function can be obtained by replacing $i\Theta(t)$ by $-i\Theta(-t)$ in Eq. (4.28). We see immediately that there is simple relation between $G_r(t)$ and $F^0(\mathbf{k},\mathbf{k}',s,s',x)$, i.e.,

$$F^0(\mathbf{k},\mathbf{k}',s,s',x)=i\frac{G_r(\mathbf{k},\mathbf{k}',s,s',x)-G_a(\mathbf{k},\mathbf{k}',s,s',x)}{e^{\beta x}-1} . \tag{4.29}$$

The equation of motion for $G_r(\mathbf{k},\mathbf{k}',s,s',x)$ is

$$\begin{aligned}\left[i\frac{\partial}{\partial t}-E_{\mathbf{k}',s'}+E_{\mathbf{k}'-\mathbf{q},s'}\right]&G_r(\mathbf{k},\mathbf{k}',s,s',t)\\
&=(n_{\mathbf{k}',s'}-n_{\mathbf{k}'-\mathbf{q},s'})\delta(\mathbf{k}+\mathbf{q}-\mathbf{k}')\delta_{s,s'}\\
&\quad-i\Theta(t)\sum_{\mathbf{k}_1,\mathbf{q}_1,s_1}V_{q_1}e_{s'}e_{s_1}\langle[a^\dagger_{\mathbf{k}'-\mathbf{q},s'}(t)a^\dagger_{\mathbf{k}_1-\mathbf{q}_1,s_1}(t)a_{\mathbf{k}_1,s_1}(t)a_{\mathbf{k}'-\mathbf{q}_1,s'}(t),a^\dagger_{\mathbf{k}+\mathbf{q},s}(0)a_{\mathbf{k},s}(0)]\rangle\\
&\quad+i\Theta(t)\sum_{\mathbf{k}_1,\mathbf{q}_1,s_1}V_{q_1}e_se_{s_1}\langle[a^\dagger_{\mathbf{k}'-\mathbf{q}+\mathbf{q}_1,s'}(t)a^\dagger_{\mathbf{k}_1-\mathbf{q}_1,s_1}(t)a_{\mathbf{k}_1,s_1}(t)a_{\mathbf{k}',s'}(t),a^\dagger_{\mathbf{k}+\mathbf{q},s}(0)a_{\mathbf{k},s}(0)]\rangle .\end{aligned} \tag{4.30}$$

This set of equations are exact. We solve it within the same approximation as was used in deriving Eq. (4.10). The result is a simplified integral equation for $G_r(\mathbf{k},\mathbf{k}',s,s',x)$,

$$\begin{aligned}[x+i\delta+E_{s'}(\mathbf{k}'-\mathbf{q})-E_{s'}(\mathbf{k}')]&G_r(\mathbf{k},\mathbf{k}',s,s',x)\\
&=\frac{[n_{s'}(\mathbf{k}'-\mathbf{q})-n_{s'}(\mathbf{k}')]}{2\pi}\delta_{s,s'}\delta_{\mathbf{k}',\mathbf{k}+\mathbf{q}}[n_{s'}(\mathbf{k}'-\mathbf{q})-n_{s'}(\mathbf{k}')]\sum_{\mathbf{k}'',s''}U_q^{s's''}e_{s''}G_r(\mathbf{k},\mathbf{k}'',s,s'',x) .\end{aligned} \tag{4.31}$$

This equation is similar to Eq. (4.14). If we compare it with Eq. (4.14), we see that the inhomogeneous term $(ie_{s'}/m_{s'})(\mathbf{q}\cdot\mathbf{E}/\epsilon)F_q^0$ in Eq. (4.14) is corresponding to the inhomogeneous term $[n_{s'}(\mathbf{k})-n_{s'}(\mathbf{k}-\mathbf{q})]\delta(\mathbf{k}+\mathbf{q}-\mathbf{k}')\delta_{s,s'}$ in Eq. (4.31). Using a similar method as we used in the solution of Eq. (4.14), and after some algebra, we obtain the solution for retarded Green's function. Similarly, the solution for advanced Green's function can be obtained. We define $x^\pm=x\pm i\delta$ where $\delta\rightarrow 0^+$, we then write our solution for $G_r=G_+$ and $G_a=G_-$ where

$$G_\pm(\mathbf{k},\mathbf{k}',e\text{-}e,x)=Q_q(\mathbf{k},x^\pm)\left[\delta(\mathbf{k}+\mathbf{q}-\mathbf{k}')+\frac{Q_q(\mathbf{k}',x^\pm)}{\tilde{\epsilon}_q(x^\pm)}\{e^2U^{e\text{-}e}[1-e^2U^{h\text{-}h}B(x^\pm)]+e^4U^{e\text{-}h}U^{h\text{-}e}B(x^\pm)\}\right] , \tag{4.32a}$$

and

$$G_\pm(\mathbf{k},\mathbf{k}',e\text{-}h,x)=\frac{Q_q(\mathbf{k},x^\pm)B_q(\mathbf{k}',x^\pm)}{\tilde{\epsilon}_q(x^\pm)}\{ee'U^{h\text{-}e}[1-e^2U^{h\text{-}h}B(x^\pm)]+e^2ee'U^{h\text{-}h}U^{h\text{-}e}B(x^\pm)\} , \tag{4.32b}$$

and

$$G_\pm(\mathbf{k},\mathbf{k}',h\text{-}h,x)=B_q(\mathbf{k},x^\pm)\left[\delta(\mathbf{k}+\mathbf{q}-\mathbf{k}')+\frac{B_q(\mathbf{k}',x^\pm)}{\tilde{\epsilon}_q(x^\pm)}\{e^2U^{h\text{-}h}[1-e^2U^{e\text{-}e}Q(x^\pm)]+e^4U^{h\text{-}e}U^{e\text{-}h}Q(x^\pm)\}\right] , \tag{4.32c}$$

and

$$G_\pm(\mathbf{k},\mathbf{k}',h\text{-}e,x)=\frac{Q_q(\mathbf{k},x^\pm)B_q(\mathbf{k}',x^\pm)}{\tilde{\epsilon}_q(x^\pm)}\{ee'U^{e\text{-}h}[1-e^2U^{e\text{-}e}Q(x^\pm)]+e^2ee'E^{e\text{-}e}U^{e\text{-}h}Q(x^\pm)\} . \tag{4.32d}$$

Here

$$Q_q(\mathbf{k},x^\pm)=\frac{n_e(\mathbf{k})-n_e(\mathbf{k}-\mathbf{q})}{x^\pm+E_e(\mathbf{k})+E_e(\mathbf{k}-\mathbf{q})} ,$$

and

$$B_q(\mathbf{k},x^{\pm})=\frac{n_h(\mathbf{k})-n_h(\mathbf{k}-\mathbf{q})}{x^{\pm}+E_h(\mathbf{k})+E_h(\mathbf{k}-\mathbf{q})} .$$

In our expression for conductivity, the quantity that needs to be determined is

$$\begin{aligned}\sum_{\mathbf{k},\mathbf{k}'}\frac{F_q^0(\mathbf{k},\mathbf{k}',e,e,x)}{x+\omega-E_e(\mathbf{k}')+E_e(\mathbf{k}'-\mathbf{q})}&=\frac{1}{e^{\beta x}-1}\sum_{\mathbf{k},\mathbf{k}'}\frac{G_r^0(\mathbf{k},\mathbf{k}',e,e,x)-G_a^0(\mathbf{k},\mathbf{k}',e,e,x)}{x+\omega-E_e(\mathbf{k}')+E_e(\mathbf{k}'-\mathbf{q})}\\&=\frac{1}{e^{\beta x}-1}[1-e^2U^{h\text{-}h}B_q(x^+)]\frac{Q_q(x^++\omega)-Q_q(x^+)}{\omega\tilde{\epsilon}_q(x^+)}\\&\quad-\frac{1}{e^{\beta x}-1}[1-e^2U^{h\text{-}h}B_q(x^-)]\frac{Q_q(x^++\omega)-Q_q(x^-)}{\omega\tilde{\epsilon}_q(x^-)} ,\end{aligned}\tag{4.33}$$

and

$$\begin{aligned}\sum_{\mathbf{k},\mathbf{k}'}\frac{F_q^0(\mathbf{k},\mathbf{k}',e,h,x)}{x+\omega-E_h(\mathbf{k}')+E_h(\mathbf{k}'-\mathbf{q})}&=\frac{1}{e^{\beta x}-1}\sum_{\mathbf{k},\mathbf{k}'}\frac{G_r^0(\mathbf{k},\mathbf{k}',e,h,x)-G_r^0(\mathbf{k},\mathbf{k}',e,h,x)}{x+\omega-E_h(\mathbf{k}')+E_h(\mathbf{k}'-\mathbf{q})}\\&=\frac{ee'U^{h\text{-}e}Q_q(x^+)}{e^{\beta x}-1}\frac{B_q(x^++\omega)-B_q(x^+)}{\omega\tilde{\epsilon}_q(x^+)}-\frac{ee'U^{h\text{-}e}Q_q(x^-)}{e^{\beta x}-1}\frac{B_q(x^++\omega)-B_q(x^-)}{\omega\tilde{\epsilon}_q(x^-)} ,\end{aligned}\tag{4.34}$$

and

$$\begin{aligned}\sum_{\mathbf{k},\mathbf{k}'}\frac{F_q^0(\mathbf{k},\mathbf{k}',h,e,x)}{x+\omega-E_e(\mathbf{k}')+E_e(\mathbf{k}'-\mathbf{q})}&=\frac{1}{e^{\beta x}-1}\sum_{\mathbf{k},\mathbf{k}'}\frac{G_r^0(\mathbf{k},\mathbf{k}',h,e,x)-G_r^0(\mathbf{k},\mathbf{k}',h,e,x)}{x+\omega-E_e(\mathbf{k}')+E_e(\mathbf{k}'-\mathbf{q})}\\&=\frac{ee'U^{e\text{-}h}B_q(x^+)}{e^{\beta x}-1}\frac{Q_q(x^++\omega)-Q_q(x^+)}{\omega\tilde{\epsilon}_q(x^+)}-\frac{ee'U^{e\text{-}h}B_q(x^-)}{e^{\beta x}-1}\frac{Q_q(x^++\omega)-Q_q(x^-)}{\omega\tilde{\epsilon}_q(x^-)} ,\end{aligned}\tag{4.34'}$$

and

$$\begin{aligned}\sum_{\mathbf{k},\mathbf{k}'}\frac{F_q^0(\mathbf{k},\mathbf{k}',h,h,x)}{x+\omega-E_h(\mathbf{k}')+E_h(\mathbf{k}'-\mathbf{q})}&=\frac{1}{e^{\beta x}-1}\sum_{\mathbf{k},\mathbf{k}'}\frac{G_r^0(\mathbf{k},\mathbf{k}',h,h,x)-G_a^0(\mathbf{k},\mathbf{k}',h,h,x)}{x+\omega-E_h(\mathbf{k}')+E_h(\mathbf{k}'-\mathbf{q})}\\&=\frac{1}{e^{\beta x}-1}[1-e^2U^{e\text{-}e}Q_q(x^+)]\frac{B_q(x^++\omega)-B_q(x^+)}{\omega\tilde{\epsilon}_q(x^+)}\\&\quad-\frac{1}{e^{\beta x}-1}[1-e^2U^{e\text{-}e}Q_q(x^-)]\frac{B_q(x^++\omega)-B_q(x^-)}{\omega\tilde{\epsilon}_q(x^-)} .\end{aligned}\tag{4.35}$$

Substituting Eqs. (4.33)–(4.35) in Eq. (4.21), we obtain

$$\begin{aligned}T_q(e\text{-}h)&=\sum_{\mathbf{k}}\tilde{T}_q(e\text{-}h)\\&=\frac{ee'U^{h\text{-}e}\mathbf{q}\cdot\mathbf{E}}{\omega^2\tilde{\epsilon}_q(\omega+x^+)}\frac{1}{\tilde{\epsilon}_q(x^+)}\Bigg\{\frac{e'}{m_h}[1-e^2U^{e\text{-}e}Q_q(x^++\omega)]Q_q(x^+)[B_q(x^++\omega)-B_q(x^+)]\\&\qquad\times\frac{e}{m_e}[1-e^2U^{h\text{-}h}B_q(x^+)]Q_q(x^+)[Q_q(x^++\omega)-Q_q(x^+)]\Bigg\}\\&\quad-\frac{1}{\tilde{\epsilon}_q(x^-)}\Bigg\{\frac{e'}{m_h}[1-e^2U^{e\text{-}e}Q_q(x^++\omega)]Q_q(x^-)[B_q(x^++\omega)-B_q(x^-)]\\&\qquad\times\frac{e}{m_e}[1-e^2U^{h\text{-}h}B_q(x^-)]Q_q(x^-)[Q_q(x^++\omega)-Q_q(x^-)]\Bigg\} ,\end{aligned}\tag{4.36}$$

and $T_q(h\text{-}e)$ can be obtained by interchanging $i\leftrightarrow e$ and $B\leftrightarrow Q$. Using the expression (4.30) we have

$$U^{e\text{-}h}T_q(e\text{-}h)-U^{h\text{-}e}T_q(h\text{-}e)$$
$$=\frac{ee'U^{h\text{-}e}\mathbf{q}\cdot\mathbf{E}}{\omega^2\tilde{\epsilon}_q(\omega+x^+)}\left[\frac{e}{me}-\frac{e'}{m_h}\right]\left[\frac{1}{\tilde{\epsilon}_q(x^+)}[Q_q(x^++\omega)-Q_q(x^+)][B_q(x^++\omega-B_q(x^+)]\right.$$
$$\left.-\frac{1}{\tilde{\epsilon}_q(x^-)}[Q_q(x^++\omega)-Q_q(x^-)][B_q(x^++\omega)-B_q(x^-)]\right] . \tag{4.37}$$

Our final complete expression for current $\langle \mathbf{J}_1\rangle$ is

$$\langle \mathbf{J}_1\rangle=\frac{e^4}{2\omega^3}\left[\frac{e}{m_e}-\frac{e'}{m_h}\right]^2\int\frac{d\mathbf{q}}{(2\pi)^3}V_qU_q^{e\text{-}h}\mathbf{q}\mathbf{q}\cdot\mathbf{E}\frac{\mathrm{P}}{2\pi i}$$
$$\times\int dx\coth\left[\frac{\beta x}{2}\right]\left[\frac{1}{\tilde{\epsilon}_q(x^++\omega)\tilde{\epsilon}_q(x^+)}[Q_q(x^++\omega)-Q_q(x^+)][B_q(x^++\omega)-B_q(x^+)]\right.$$
$$\left.-\frac{1}{\tilde{\epsilon}_q(x^++\omega)\tilde{\epsilon}_q(x^-)}[Q_q(x^++\omega)-Q_q(x^-)][B_q(x^++\omega)-B_q(x^-)]\right] . \tag{4.38}$$

This is a rather complicated result but in principle can be evaluated numerically for specific problems. Besides the three-dimensional integration in Eq. (4.38), we need the numerical solution for $U^{ss'}$. Our result is temperature dependent and thus is applicable for quantum plasmas as well as for nondegenerate classical plasmas provided the plasma parameter r_s is of order of unity. Before presenting numerical computation, we would like to consider our result in two limiting cases.

a. High-temperature limit. In order to obtain the nondegenerate or classical case, we must take the high-temperature limit of Eq. (4.38). We first make the transformation $x=huq$ and obtain

$$Q_q^{\pm}(u,s)=\frac{n_s}{m_s}\int d\mathbf{v}\frac{\mathbf{q}\cdot\dfrac{\partial f_0(\mathbf{v},s)}{\partial\mathbf{v}}}{\mathbf{q}\cdot\mathbf{v}-qu\mp i\alpha}=-n_s\beta[1+uf_s^{\pm}(u)] , \tag{4.39}$$

where $Q_q(u,s)$ stands for $Q_q(u)$ or $B_q(u)$. In Eq. (4.39) use has been made of the following transformations:

$$\mathbf{k}=\frac{m_s\mathbf{v}}{\hbar} ,$$
$$n_s(\mathbf{k}+\mathbf{q})-n_s(\mathbf{k})=\left[\frac{2\pi\hbar}{m_s}\right]^3\frac{n_s\hbar}{m_s}\mathbf{q}\cdot\frac{\partial f_0(\mathbf{v},s)}{\partial\mathbf{v}} , \tag{4.40}$$

where $f_0(\mathbf{v},s)$ is the Maxwell distribution for the s species, and

$$f_s^{\pm}(u)=\int du'\frac{f_s(u')}{u'-u\mp i\alpha} , \tag{4.41}$$

with $f_s(u)$ being the one-dimensional Maxwell-Boltzmann distribution function. In the high-temperature limit

$$\coth\left[\frac{\beta x}{2}\right]\approx\frac{1}{\beta q\hbar u} , \tag{4.42}$$

and thus

$$\langle \mathbf{J}_1\rangle=\frac{e^4n^2\beta}{2\omega^3}\left[\frac{e}{m_e}-\frac{e'}{m_\hbar}\right]^2\int_0^{q_{\max}}d\mathbf{q}\frac{1}{(2\pi)^3}V_qU_q^{e\text{-}h}\mathbf{q}\mathbf{q}\cdot\mathbf{E}\frac{\mathrm{P}}{2\pi i}$$
$$\times\int\frac{du}{u}\left[\frac{[(u+w)f_e^+(u+w)-uf_e^+(u)][(u+w)f_h^+(u+w)-uf_h^+(u)]}{\tilde{\epsilon}_q(u+w)\tilde{\epsilon}_q(u)}\right.$$
$$\left.-\frac{[(u+w)f_e^+(u+w)-uf_e^-(u)][(u+w)f_h^+(u+w)-uf_h^-(u)]}{\tilde{\epsilon}_q(u+w)\tilde{\epsilon}^*_q(u)}\right] , \tag{4.43}$$

where $w=\omega/q$ and $q_{\max}$ is the cutoff due to large angle collision (in the quantum case $q_{\max}$ is given by the recoil). One should note here that for the nondegenerate case $U_q^{ss'}$ (and thus $\tilde{\epsilon}$) must be obtained in the high-frequency limit. Thus Eq. (4.43) is our generalization of the result of Ref. 3 which includes the effect of short-range correlations between the charged particles.

b. Heavy-hole limit. In order to carry out the limit of fixed-ion scatterers we treat the ion (hole) classically, namely, we replace the Q's by their classical representation Eq. (4.40), and prescribe for the ions

$$f_h^{\pm}(u) = -\mathrm{P}\frac{1}{u} \pm i\pi\delta(u) \ . \tag{4.44}$$

One should notice that in this limit

$$\tilde{\epsilon}(q,x) \rightarrow 1 + U_q^{e\text{-}e} Q_q(x) \ , \tag{4.45}$$

while

$$\coth\left[\frac{\beta x}{2}\right][B(x+\omega) \pm B(x)] \rightarrow \pm i n \delta(x) \ .$$

Thus we obtain

$$\langle \mathbf{J}_1 \rangle = \frac{2e^4 n}{2\omega^3}\frac{e^2}{m_e^2}\int \frac{d\mathbf{q}}{(2\pi)^3} V_q U_q^{e\text{-}h} \mathbf{q}\mathbf{q}\cdot\mathbf{E}\frac{Q_q(\omega)-Q_q(0)}{\tilde{\epsilon}_q(\omega)\tilde{\epsilon}_q(0)}$$
$$= \frac{2e^4 n}{2\omega^3}\frac{e^2}{m_e^2}\int \frac{d\mathbf{q}}{(2\pi)^3}\frac{V_q U_q^{e\text{-}h}}{U_q^{e\text{-}e}} \times \mathbf{q}\mathbf{q}\cdot\mathbf{E}\left[\frac{1}{\tilde{\epsilon}_q(\omega)} - \frac{1}{\tilde{\epsilon}_q(0)}\right] \ . \tag{4.46}$$

For the one-component plasma $\tilde{\epsilon}$ can be related to the dielectric function through

$$\frac{V_q Q_q}{\tilde{\epsilon}} = \frac{1}{\epsilon} - 1$$

where $\epsilon_q(\omega)$ is the dielectric function for correlated plasma first introduced by Singwi-Tosi-Land-Sjolander (Ref. 7), i.e.,

$$\epsilon_q(\omega) = 1 + \frac{V_q Q_q(\omega)}{1 - W^{e\text{-}e} V_q Q_q(\omega)} \ . \tag{4.47}$$

The quantity $W^{e\text{-}e}$ is the local field correction defined as $U_q^{e\text{-}e}/V_q = 1 + W^{e\text{-}e}$. Thus Eq. (4.46) is our generalization of the electron-ion system[1-5] which includes the short-range correlations.

V. RESISTIVITY AND COLLISION FREQUENCY

Let us consider a system consisting of equal number of electrons and holes. The total current can be written as

$$\langle \mathbf{J}_1 \rangle = \frac{ine^2\mathbf{E}}{\omega}\left[\frac{e}{m_e} - \frac{e'}{m_h}\right] + \frac{\mathbf{E}e^2}{3\omega^3}\left[\frac{e}{m_e} - \frac{e'}{m_h}\right]^2$$
$$\times \int \frac{dq\, q^4}{2\pi^2} V_q U_q^{e\text{-}h} \frac{\mathrm{P}}{2\pi i}\int dx \coth\left[\frac{\beta x}{2}\right]\left[\frac{1}{\tilde{\epsilon}_q(x^+ + \omega)\tilde{\epsilon}_q(x^+)}[Q_q(x^+ + \omega) - Q_q(x^+)][B_q(x^+ + \omega) - B_q(x^+)]\right.$$
$$- \frac{1}{\tilde{\epsilon}_q(x^+ + \omega)\tilde{\epsilon}_q(x^-)}[Q_q(x^+ + \omega) - Q_q(x^-)]$$
$$\left.\times [B_q(x^+ + \omega) - B_q(x^-)]\right] \ , \tag{5.1}$$

and the conductivity can be written as

$$\sigma(\omega) = \frac{ine^2}{\mu\omega}\left[1 + \frac{I(\omega)}{\omega}\right] = \sigma_0(\omega)\left[1 + \frac{I(\omega)}{\omega}\right] \ , \tag{5.2}$$

where

$$\frac{1}{\mu} = \frac{1}{m_e} + \frac{1}{m_h}$$

and

$$I(\omega) = \frac{e^4}{3n\omega\mu}\int dq \frac{q^4}{2\pi^2} V_q U_q^{e\text{-}h}\frac{\mathrm{P}}{2\pi i}$$
$$\times \int dx \coth\left[\frac{\beta x}{2}\right]\left[\frac{1}{\tilde{\epsilon}_q(x^+ + \omega)\tilde{\epsilon}_q(x^+)}[Q_q(x^+ + \omega) - Q_q(x^+)][B_q(x^+ + \omega) - B_q(x^+)]\right.$$
$$\left. - \frac{1}{\tilde{\epsilon}_q(x^+ + \omega)\tilde{\epsilon}_q(x^-)}[Q_q(x^+ + \omega) - Q_q(x^-)][B_q(x^+ + \omega) - B_q(x^-)]\right] \ . \tag{5.3}$$

Compare Eqs. (5.2) and (5.3) with the Drude formula for conductivity:

$$\sigma(\omega)=\frac{ine^2}{\mu(\omega+i\nu)} \ . \tag{5.4}$$

We obtain for high frequency ($\omega \gg \nu$)

$$\nu=\text{Im}[I(\omega)]=\nu_0+\nu_1 \ , \tag{5.5}$$

where ν is the collision frequency. ν_0 is the collision frequency calculated from RPA and ν_1 is the correction on collision frequency due to the short-range correlation. Taking the imaginary part of $I(\omega)$ and using the analytical properties of Q, B, and "dielectric function," respectively, we can write

$$\nu=\frac{\pi e^4}{3\omega n\mu}\int dq\frac{q^4}{2\pi^2}V_q U_q^{e\text{-}h}\frac{\text{P}}{2}\int dx\left[\coth\left[\frac{\beta x}{2}\right]-\coth\left[\frac{\beta(x+\omega)}{2}\right]\right]F(x,x+\omega) \ , \tag{5.6}$$

where

$$F(x,x+\omega)=\frac{F_1+F_2+F_3}{|\tilde{\epsilon}(x)|^2\,|\tilde{\epsilon}(x+\omega)|^2} \ , \tag{5.7}$$

with

$$F_1=[B_2(x+\omega)+B_2(x)]\{[\tilde{\epsilon}_1(x+\omega)\tilde{\epsilon}_1(x)+\tilde{\epsilon}_2(x+\omega)\tilde{\epsilon}_2(x)][Q_2(x+\omega)+Q_2(x)]$$
$$-[\tilde{\epsilon}_1(x)\tilde{\epsilon}_2(x+\omega)-\tilde{\epsilon}_1(x+\omega)\tilde{\epsilon}_2(x)][Q_1(x+\omega)-Q_1(x)]\} \ , \tag{5.8}$$

and

$$F_2=[B_2(x)-B_2(x+\omega)]\{[\tilde{\epsilon}_1(x+\omega)\tilde{\epsilon}_1(x)-\tilde{\epsilon}_2(x+\omega)\tilde{\epsilon}_2(x)][Q_2(x+\omega)-Q_2(x)]$$
$$-[\tilde{\epsilon}_1(x)\tilde{\epsilon}_2(x+\omega)+\tilde{\epsilon}_1(x+\omega)\tilde{\epsilon}_2(x)][Q_1(x+\omega)-Q_1(x)]\} \ , \tag{5.9}$$

and

$$F_3=2[B_1(x+\omega)-B_1(x)]\{\tilde{\epsilon}_1(x+\omega)\tilde{\epsilon}_2(x)Q_2(x+\omega)-\tilde{\epsilon}_1(x)\tilde{\epsilon}_2(x+\omega)Q_2(x)-[Q_1(x+\omega)-Q_1(x)]\tilde{\epsilon}_2(x)\tilde{\epsilon}_2(x+\omega)]\} \ , \tag{5.10}$$

where $\tilde{\epsilon}_1$ and $\tilde{\epsilon}_2$ represent, respectively, the real and imaginary parts of the dielectric function. Similarly, Q_1 and Q_2 (B_1 and B_2) represent, respectively, the real and the imaginary parts of $Q(B)$. At $T=0$, the factor $\{\coth(\beta x/2)-\coth[\beta(x-\omega)/2]\}$ reduces to 2 within the region $-\omega<x<0$ and vanishes outside. Therefore Eq. (5.6) can be written as

$$\nu=\frac{\pi e^4}{3\omega n\mu^2}\int dq\frac{q^4}{2\pi^2}V_q U_q^{e\text{-}h}\int_0^{\omega}dx\,F(x,\omega-x) \ . \tag{5.11}$$

We shall evaluate this equation for some typical parameters and discuss the effect of correlation in Sec. VI.

VI. DISCUSSION

We have derived an expression for the conductivity and collision frequency for a two-component plasma including short-range correlation. We found that the short-range correlation will affect the dynamical conductivity in two ways. First, we found that the Coulomb matrix element V_q is replaced by the effective interaction $U_q^{e\text{-}h}$ in Eq. (5.6). Note that only one of the matrix elements V_q is replaced by the effective interaction (due to the short-range correlations). The other Coulomb matrix element as it appears in Eq. (3.10) is part of the expression for $\langle \mathbf{J}_1\rangle$ and is not part of the density-density correlation function. This replacement indicating the change of the scattering matrix of the electrons by the holes which in our theory are taken to be correlated rather than described by plane waves. This in turn enhances the collision frequency. Second, the correlation affects the dielectric function. Here $U^{e\text{-}e}$, $U^{h\text{-}h}$, and $U^{e\text{-}h}$ all contribute to the change of screening effect. Here the replacement of V_q by appropriate effective potential in the dielectric function tends to reduce the screening effect. We point out that the short-range correlation breaks the symmetry between the magnitude of the electron-electron interaction versus electron-hole interaction ($U_q^{e\text{-}e}\neq U_q^{e\text{-}h}$). Therefore, the dielectric function includes an extra term which is proportional to the product of density fluctuations of electron and hole (QB) as can be seen in Eq. (4.22). This extra term will further reduce the effect of screening. The combined effect of enhanced scattering and reduced screening will increase collision frequency or the absorption constant.

We have performed some numerical calculations of the collision frequency at zero temperature. We use the numerical values for the local field correction $W^{e\text{-}e}$, $W^{h\text{-}h}$, and $W^{e\text{-}h}$ given by Vashishta *et al.*[19] The results for collision frequency are shown in Figs. 1–3.

In Fig. 1 we consider electron-hole plasma with equal densities of electrons and holes, for a plasma parameter

$r_s=1$. The two dashed curves represent the collision frequency for mass ratio $\alpha=1$ and 4, without including short-range correlations. The solid curves represent our solution for the collision frequency including short-range correlations for the same system. Similarly, in Fig. 2, we plot the collision frequency for electron-hole plasma for $r_s=2$. Here the results for three values of mass ratio are presented, i.e., $\alpha=1,4,6$. As before, the dashed curve represents the collision frequency when short-range correlations are omitted and the solid curve gives the collision frequency when short-range correlations are taken into account.

We point out that for a one-component plasma, RPA overestimates the screening effect at large wave numbers. Here the short-range correlation tends to decrease the screening at short distances (large q) due to the particles' repulsion. For two-component plasma the situation is more complicated. For example, at short distances, an electron will experience less screening by other electrons but an enhanced screening by the holes. No physical arguments can determine the effect of short-range correlations on the screening without detailed calculations. We found that for the values of r_s used in our paper the screening was less effective than what is predicted in RPA. In this paper we calculate the collision rate for long-wavelength radiation fields due to electron-hole scattering. In our case, less screening results in more efficient scattering and the increase of ν. Moreover, our theory takes into account the attractive electron-hole correlations during the scattering process. We calculate the electron-hole scattering matrix including correlations. We take into account the increase of the electron density around the hole during the collision process. It is worthwhile to mention that the electron-hole correlations will eventually, at large enough values of r_s, result in the formation of an exciton gas. Our results show an increase in the collision frequency due to short-range correlations for the values of r_s and mass ratio considered in this paper. However, to better understand the behavior of the collision frequency as a function of r_s and α, more numerical work is needed. We at present are limited by the numerical solutions of the effective interactions for the parameters presented in Ref. 19.

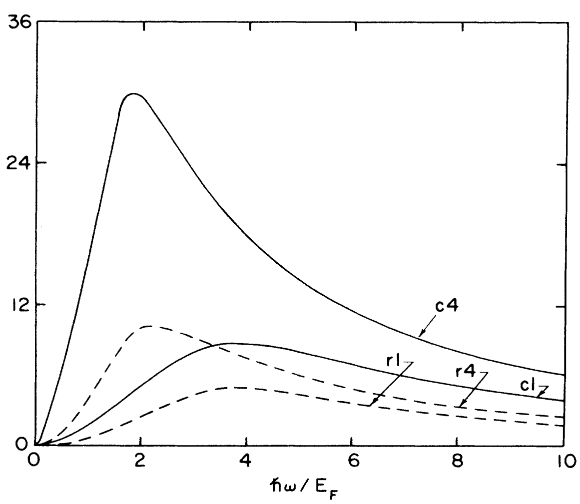

FIG. 1. Plot of normalized collision frequency $\tilde{\nu}=20\pi\nu/E_F(e)$ as a function of normalized frequency ω/E_F. $r_s=1.0$; $C1$, present theory with $\alpha=1$; $C4$, present theory with $\alpha=4.0$. $r1$, RPA theory with $\alpha=1$; $r4$, RPA theory with $\alpha=4.0$.

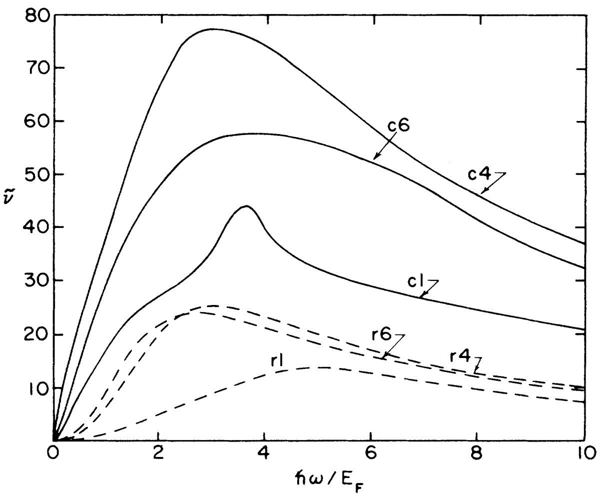

FIG. 2. Plot of normalized collision frequency $\tilde{\nu}=20\pi\nu/E_F(e)$ as a function of normalized frequency ω/E_F. $r_s=2.0$; $C1$, present theory with $\alpha=1$; $C4$, present theory with $\alpha=4.0$; $C6$, present theory with $\alpha=6.0$. $r1$, RPA theory with $\alpha=1$; $r4$, RPA theory with $\alpha=4.0$.

From Fig. 3, we find that the effect of short-range correlations at low frequency is much more important than that at high frequency. The large correction at low frequency is due to the change of the screening. At low frequency, $\tilde{\epsilon}(q,\omega)$ can be approximated by its static value $\tilde{\epsilon}(q)$. The collision frequency ν can be given as

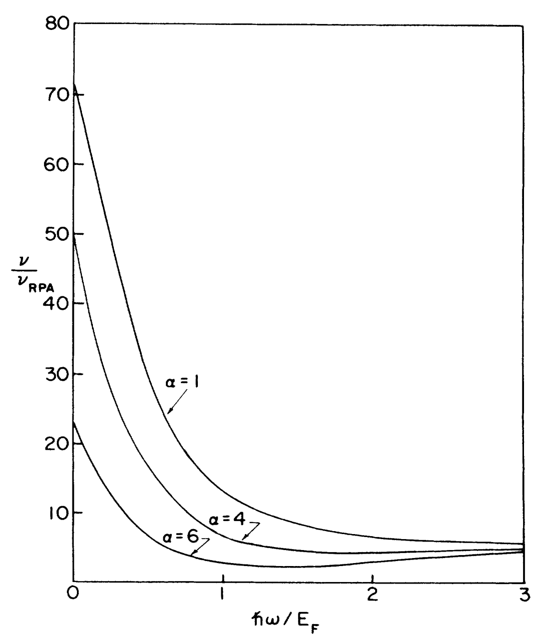

FIG. 3. Plot of the ratio of the collision frequencies between present theory and RPA for $r_s=2.0$ and $\alpha=1,4,6$.

$$\nu = A\omega^2 \int_0^1 dq \frac{1}{q^2} \frac{1+W^{e\text{-}h}}{[\tilde{\epsilon}(q)]^2} , \tag{6.1}$$

where A is a constant given by $A = (\pi/24) r_s (1+\alpha)$. The quantitative enhancement factor for the collision frequency due to the reduction of screening and the enhancement of the scattering matrix element $(1+W^{e\text{-}h})$ at low frequencies can be 10 to 15 times larger compared to that calculated from RPA. We also found, as expected, that the effect of correlation on the scattering matrix is larger for large r_s. However, the effect of correlation due to screening is more important at small r_s, i.e., for high densities, and when r_s increases the effect of the screening is reduced.

In conclusion, we have calculated the dynamical conductivity and the collision frequency in a two-component plasma. The short-range correlation is taken into account. Numerical results showing quantitative effect of correlation are presented.

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