Finite temperature effects in light scattering off Cooper-paired Fermi atoms

Bimalendu Deb

Department of Materials Science and Centre for Atomic, Molecular and Optical Sciences, Indian Association for the Cultivation of Science, Jadavpur, Kolkata-700 032, India

Received 3 April 2007, in final form 11 May 2007
Published 5 June 2007
Online at stacks.iop.org/JPhysB/40/2399

Abstract

We study stimulated light scattering off a superfluid Fermi gas of atoms at finite temperature. We derive a response function that takes into account vertex correction due to final state interactions and analyse finite temperature effects on collective and quasiparticle excitations of a uniform superfluid Fermi gas. Light polarization is shown to play an important role in excitations. Our results suggest that it is possible to excite Bogoliubov–Anderson phonons at large scattering lengths by light scattering.

(Some figures in this article are in colour only in the electronic version)

1. Introduction

Since the first realization of Bose–Einstein condensation in atomic gases in 1995 [1], there has been tremendous growth in research activities with cold atoms. Recent experimental studies [2–8] with cold fermionic atoms have generated renewed interest in quantum many-body physics. Atomic Fermi gases in traps provide a unique laboratory system for exploring the physics of interacting fermions with tunable interactions. Fermi degeneracy in a trapped atomic gas was first demonstrated by DeMarco and Jin in 1999 [9]. In the recent past, there has been many reports of possible observations of Fermi superfluidity (FS). However, unambiguous evidence of FS was demonstrated by Zwierlein et al [10]. The detection of pairing gap [11–15] and collective modes [16–18] of FS are current issues of interest. The physics at the crossover [19–21] between BCS state of atoms and BEC of molecules [22] formed from Fermi atoms is of prime interest. A number of theoretical investigations [23] have dealt with FS near the crossover. Two very recent experiments [24, 25] on two-component Fermi gases with imbalanced spin components provide new insight on the nature of FS and perhaps point to some new state of matter known as interior gap (IG) superfluidity predicted by Liu and Wilczek [26]. The occurrence of IG superfluidity in a two-component Fermi gas was theoretically predicted by Deb et al [27].
In order to study the nature of FS, it is important to derive the appropriate response function of Cooper-paired Fermi atoms due to an external perturbation (such as photon or rf field). Our purpose here is to calculate response function of superfluid Fermi gas at finite temperature due to stimulated light scattering. In a previous paper [28], we derived the response function at zero temperature. We have also shown that it is possible to excite selectively a single partner atom (of a particular hyperfine spin state) of a Cooper-pair exploiting light polarizations in the presence of a strong magnetic field. We present here a detailed method of calculation of response function at finite temperature due to light scattering. We study the effects of finite temperature and light polarization on single-particle excitation as well as the collective mode of density fluctuations. This collective mode known as Bogoliubov–Anderson (BA) phonon [29] has been theoretically studied in the context of fermionic atoms [30]. We show here that it may be possible to excite this mode by light scattering: This mode appears as a resonance in dynamic structure versus energy transfer in scattering. At finite temperature, the resonance becomes broadened due to Landau damping. In case of single-particle excitation spectrum when Cooper-pairs are broken \((\omega > 2\Delta)\), the peak occurs at a higher energy as the temperature is lowered. If the momentum transfer is higher, the single-particle spectrum becomes sharper. Light polarization significantly affects the excitation spectrum. Circular polarization of light leads to a positive shift of the peak of the single-particle spectrum at lower temperatures.

This paper is organized in the following way. In the following section, we discuss in brief the stimulated light scattering off Cooper-paired Fermi atoms. In section 3, we present in detail the derivation of the response function with vertex correction at finite temperature. We next describe numerical results in section 4. The paper is concluded in section 5.

2. Stimulated photon scattering

Stimulated light scattering will occur when two laser beams with nearly equal frequencies are impinged on atomic gas and tuned far-off the resonance of a transition frequency of the atoms. Let us specifically consider the two-component Fermi gas of $^6\text{Li}$. Here the two components imply the two hyperfine spin states of $F = 1/2$ ground level. As discussed previously [28], large Zeeman shifts of hyperfine sub-levels near Feshbach resonance allow us to utilize light polarization to control scattering to a significant extent. Figure 1 of [28] describes the polarization-selective light scattering by which the amount of momentum and energy transfer to the two partner atoms of a Cooper-pair can be controlled. It is thereby possible to scatter photons from either spin component. However because of long-range correlation between two components due to s-wave Cooper-pairing, the other spin component will also be affected by photon scattering. When light fields are treated classically, the effective atom–field interaction Hamiltonian is

$$H_I \propto \sum_{\sigma, k} \gamma_{\sigma, \sigma} a_{\sigma, k+q}^\dagger a_{\sigma, k}$$

where $a_{\sigma, k}(a_{\sigma, k}^\dagger)$ represents the annihilation (creation) operator of an atom with hyperfine spin $\sigma$ and center-of-mass momentum $k$, $q$ is the momentum transfer and $\gamma_{\sigma, \sigma}$ is the bare vertex corresponding scattering without any change in the spin state $\sigma$.

To describe density–density correlation function, we define the density operators by

$$\rho_{\sigma}^{(0)} = \sum_{\sigma, k} a_{\sigma, k+q}^\dagger a_{\sigma, k}$$

and

$$\rho_{\sigma}^{(\gamma)} = \sum_{\sigma, k} \gamma_{\sigma, \sigma} a_{\sigma, k+q}^\dagger a_{\sigma, k}.\quad (2)$$
3. Response function

The response function due to the density fluctuation is

$$\chi(q, \tau - \tau') = -\langle \tau \rangle [\rho_q^{\langle 0 \rangle}(\tau) \rho_q^{\langle 0 \rangle}(\tau')]$$, \hspace{1cm} (3)

where $\langle \cdots \rangle$ implies thermal averaging and $\tau_i$ is the complex $\tau$ ordering operator. By generalized fluctuation–dissipation theorem, the dynamic structure factor is given by

$$S(q, \omega) = -\frac{1}{\pi} [1 + n_B(\omega)] \text{Im}[\chi(q, z = \omega + i\delta)],$$ \hspace{1cm} (4)

where $\chi(q, z)$ represents the Fourier transform of $\chi(q, \tau)$.

3.1. Vertex equation

Within the framework of Nambu–Gorkov formalism of superconductivity, using Pauli matrices $\tau_i$, the vertex function can be written as [31]

$$\Gamma(k_+, k_-) = \tilde{\gamma} - \int \beta^{-1} \frac{d^3k}{(2\pi)^3} \sum_n \tau_3 G(k'_+) \Gamma(k'_+, k'_-) G(k'_-) \tau_3 V(k, k'),$$ \hspace{1cm} (5)

where $k_+ \equiv \{k + q/2, i(\omega_n + \nu_m/2)\}$ and $k_- \equiv \{k - q/2, i(\omega_n - \nu_m/2)\}$ with $\omega_n/\hbar = (2n + 1)(\hbar\beta)^{-1}$ and $\nu_m/\hbar = 2m(\hbar\beta)^{-1}$ representing the fermionic and bosonic Matsubara frequency, respectively. The Green’s function can be expressed in a matrix form as

$$G(k) = -\frac{i\mu_n \tau_0 + \xi_k \tau_3 + \Delta_k \tau_1}{\mu_n^2 + E_k^2},$$ \hspace{1cm} (6)

where $\mu_n/\hbar$ is the Matsubara frequency, $E_k = \sqrt{\xi_k^2 + \Delta_k^2}$ and $\xi_k = \epsilon_k - \mu$ with $\epsilon_k = \hbar^2 k^2/(2m)$. The bare vertex can be written in a matrix form

$$\tilde{\gamma} = \gamma_0 \tau_0 + \gamma_3 \tau_3,$$ \hspace{1cm} (7)

where $\gamma_0 = [\gamma_{\uparrow \uparrow} - \gamma_{\downarrow \downarrow}]/2$ and $\gamma_3 = [\gamma_{\uparrow \downarrow} + \gamma_{\downarrow \uparrow}]/2$.

The response function is related to the vertex function by

$$\chi(q, \omega) = \int \frac{d^3k}{(2\pi)^3} \beta^{-1} \sum_n \text{Tr} [\tilde{\gamma} G(k_n) \Gamma(k_n, k-) G(k_-)].$$ \hspace{1cm} (8)

To solve the vertex equation, let us expand the vertex function in terms of Pauli matrices as

$$\Gamma(k_+, k_-) = \sum_{i=0}^{3} \Gamma^{(i)}(k, q, i\nu_m) \tau_i.$$ \hspace{1cm} (9)

We replace $V(k, k')$ by an effective mean-field potential $V_{\text{eff}}$. In the weak-coupling limit, this potential reduces to the form $V_{\text{weak}} = g a_s$, where $g = 4\pi \hbar^2/m$ and $a_s$ is the s-wave scattering length. Since we are interested only in the Cooper-pairing regime, we consider the attractive interaction only and hence $a_s$ is assumed to be negative. On replacing $V(k, k')$ by $g a_s$, the gap $\Delta$ becomes $k$-independent. The vertex function also becomes $k$-independent but remains a function of $q$ and $i\nu_m$ only.

Using equations (6) and (9) in equation (5), we can write

$$\Gamma(k, q, i\nu_m) = \tilde{\gamma}_k - g a_s \int \frac{d^3k'}{(2\pi)^3} \times \beta^{-1} \sum_n \frac{1}{[(\omega_n + \nu_m/2)^2 + E_{k'}^2][(\omega_n - \nu_m/2)^2 + E_{k'}^2]} \sum_{i=0}^{3} y_i \tau_i,$$ \hspace{1cm} (10)
where $\xi_k = h^2 k^2 / (2m) - \mu$ and

$$y_0 = \left[ i(\omega_n + v_m/2)\xi_{k'} + i(\omega_n - v_m/2)\xi_{k''} \right] \Gamma^{(3)} - i\Delta[\xi_{k'} + \xi_{k''}] \Gamma^{(2)} + \Delta[i(\omega_n + v_m/2) + i(\omega_n - v_m/2)] \Gamma^{(1)} + \left[ \xi_{k'} \xi_{k''} + (\omega_n + v_m/2)(\omega_n - v_m/2) + \Delta^2 \right] \Gamma^{(0)},$$

$$y_1 = -i[i(\omega_n + v_m/2)\xi_{k'} - i(\omega_n - v_m/2)\xi_{k''}] \Gamma^{(2)} - \Delta[i(\omega_n + v_m/2) + i(\omega_n - v_m/2)] \Gamma^{(0)} + \left[ \xi_{k'} \xi_{k''} + (\omega_n + v_m/2)(\omega_n - v_m/2) + \Delta^2 \right] \Gamma^{(1)},$$

$$y_2 = -i\Delta[i(\omega_n + v_m/2) - i(\omega_n - v_m/2)] \Gamma^{(3)} - i\Delta[\xi_{k'} + \xi_{k''}] \Gamma^{(2)} - i[i(\omega_n + v_m/2) - i(\omega_n - v_m/2)] \Gamma^{(1)} + \left[ \xi_{k'} \xi_{k''} + (\omega_n + v_m/2)(\omega_n - v_m/2) + \Delta^2 \right] \Gamma^{(0)},$$

$$y_3 = \Delta[\xi_{k'} + \xi_{k''}] \Gamma^{(1)} - i\Delta[i(\omega_n + v_m/2) - i(\omega_n - v_m/2)] \Gamma^{(2)} + \left[ \xi_{k'} \xi_{k''} - (\omega_n + v_m/2)(\omega_n - v_m/2) - \Delta^2 \right] \Gamma^{(3)} + \left[ \xi_{k'} i(\omega_n - v_m/2) + i(\omega_n + v_m/2) \xi_{k''} \right] \Gamma^{(0)}.$$

There are basically two types of Matsubara frequency sums:

$$I_1(k', q) = \beta^{-1} \sum_n \left[ \frac{\omega_n + v_m/2}{(\omega_n + v_m/2)^2 + E_{k'}^2} \right] \left[ \frac{\omega_n - v_m/2}{(\omega_n - v_m/2)^2 + E_{k'}^2} \right],$$

$$I_2(k', q) = \beta^{-1} \sum_n \frac{1}{(\omega_n + v_m/2)^2 + E_{k'}^2} \times \frac{1}{(\omega_n - v_m/2)^2 + E_{k'}^2}.$$

Now, the term such as $\omega / (\omega^2 + E_k^2)$ can be written in a separable form as

$$\frac{\omega}{\omega^2 + E_k^2} = -\frac{1}{2\imath} \left[ \frac{1}{\imath \omega + E_k} + \frac{1}{\imath \omega - E_k} \right].$$

while the term such as $1/(\omega^2 + E_k^2)$ can be decomposed as

$$\frac{1}{\omega^2 + E_k^2} = \frac{1}{2E_k} \left[ \frac{1}{\imath \omega + E_k} + \frac{1}{\imath \omega - E_k} \right].$$

The terms which are odd in frequency will not contribute to the sum and so can be omitted. Using these decompositions and after some algebra as is shown in appendix A, we obtain

$$I_1(k', q) = -\frac{1}{4} [T_{11} + T_{12} + T_{21} + T_{22}]$$

$$I_2(k', q) = \frac{1}{4E_{k'} E_{k''}} [T_{11} - T_{12} - T_{21} + T_{22}],$$

where

$$T_{11} = \frac{\tanh(\beta E_{k'}/2) - \tanh(\beta E_{k''}/2)}{2(E_{k'} - E_{k''} + iv_m)},$$

$$T_{12} = \frac{\tanh(\beta E_{k''}/2) - \tanh(\beta E_{k'}/2)}{2(E_{k'} - E_{k''} + iv_m)},$$

$$T_{21} = \frac{\tanh(\beta E_{k'}/2) - \tanh(\beta E_{k''}/2)}{2(E_{k'} - E_{k''} + iv_m)},$$

$$T_{22} = \frac{\tanh(\beta E_{k''}/2) - \tanh(\beta E_{k'}/2)}{2(E_{k'} - E_{k''} + iv_m)}.$$
\[
T_{12} = -\frac{\tanh(\beta E_{k_1}/2) + \tanh(\beta E_{k_2}/2)}{2(E_{k_1} + E_{k_2} + i\nu_m)}
\]
\[
T_{21} = -\frac{\tanh(\beta E_{k_1}/2) + \tanh(\beta E_{k_2}/2)}{2(E_{k_1} + E_{k_2} - i\nu_m)}
\]
\[
T_{22} = \frac{\tanh(\beta E_{k_1}/2) - \tanh(\beta E_{k_2}/2)}{2(E_{k_1} - E_{k_2} - i\nu_m)}
\]

In what follows, we use as the unit of energy the Fermi energy \(\epsilon_F = \hbar^2 k_F^2/(2m)\) and accordingly scale all the quantities. We denote \(\Delta = \Delta/\epsilon_F, \bar{\xi}_k = \xi_k/\epsilon_F, \tilde{k} = k/k_F\) and so on. Let \(x = \tilde{q} k \cos \theta\), where \(\theta\) is the angle between \(\mathbf{k}\) and \(\mathbf{q}\). For notational convenience, let \(E = \bar{\xi}_k + q^2/4, E_1 = E_{k_1} = \sqrt{(E - x)^2 + \Delta^2}\) and \(E_2 = E_{k_2} = \sqrt{(E + x)^2 + \Delta^2}\). Having now performed the Matsubara frequency sum, omitting the terms which are odd in \(x\), we can express the vertex equation as

\[
\Gamma(k, q, \iota, \nu_m) = \gamma_i - \kappa_s \int \frac{d^3k'}{(2\pi)^3} \sum_{i=0}^3 c_i(k')\tau_i,
\]

where \(\kappa_s = ga_s k_F^3/\epsilon_F\) and

\[
c_0 = [(\xi^2 - x^2 + \tilde{\Delta}^2)I_2 - I_1]\Gamma^{(0)},
\]
\[
c_1 = [\nu_m \Gamma^{(2)} - 2\tilde{\Delta} \Gamma^{(3)}]E I_2 + [(\xi^2 - x^2 + \tilde{\Delta}^2)I_2 + I_1]\Gamma^{(1)},
\]
\[
c_2 = \tilde{\Delta} \nu_m I_2 \Gamma^{(1)} - \nu_m E I_2 [\Gamma^{(1)} + [(\xi^2 - x^2 + \tilde{\Delta}^2)I_2 + I_1]\Gamma^{(2)},
\]
\[
c_3 = 2\tilde{\Delta} E I_3 \Gamma^{(1)} + \tilde{\Delta} \nu_m I_3 \Gamma^{(2)} + [(\xi^2 - x^2 + \tilde{\Delta}^2)I_2 - I_1]\Gamma^{(3)}.
\]

Now the vertex terms \(\Gamma^{(i)}\) form four coupled algebraic equations

\[
\Gamma^{(i)} = \gamma_i - \kappa_s \int \frac{d^3k'}{(2\pi)^3} c_i, \quad i = 0, 3
\]
\[
\Gamma^{(j)} = -\kappa_s \int \frac{d^3k'}{(2\pi)^3} c_j, \quad j = 1, 2
\]

In the limit \(q \rightarrow 0, \nu_m \rightarrow 0\) and \(\Gamma^{(2)} \rightarrow \tilde{\Delta}\), the equation for \(\Gamma^{(2)}\) reduces to the standard BCS gap equation

\[
\tilde{\Delta} = -\kappa_s \int \frac{d^3k}{(2\pi)^3} \frac{\tanh(\beta E_k/2)}{2E_k} \tilde{\Delta},
\]

which implies \(|\kappa_s|I_0 = 1\) where

\[
I_0 = \int \frac{d^3k}{(2\pi)^3} \frac{\tanh(\beta E_k/2)}{2E_k}.
\]

The gap equation expressed in this form has a logarithmic divergence. However, this divergence can be removed by renormalizing the mean-field interaction via subtracting the zero temperature and zero pairing field (\(\Delta = 0\)) part of the integral. Although this gap equation resembles the standard weak-coupling BCS gap equation, the chemical potential \(\mu\) can deviate significantly from its weak-coupling value \(\mu \simeq \epsilon_F\). The chemical potential is given by single-spin superfluid number density given by

\[
n = \frac{1}{2} \int \frac{d^3k}{(2\pi)^3} \left( 1 - \frac{\bar{\xi}_k}{E_k} \tanh(\beta E_k/2) \right).
\]
Let us first consider the off-diagonal terms $\Gamma^{(1)}$ and $\Gamma^{(2)}$. Making use of the relation (29) and the analytic continuation $i\nu_m \rightarrow \omega + i0^+$,

$$\Gamma^{(1)} = \frac{i\tilde{\omega} J_1 \Gamma^{(2)} + 2\tilde{\Delta} J_1 \Gamma^{(3)}}{(I_2 - \Delta^2 J_2 + I_1) - I_0},$$

$$\Gamma^{(2)} = \frac{i\tilde{\Delta}\tilde{\omega} J_2 \Gamma^{(3)} - i\tilde{\omega} J_1 \Gamma^{(1)}}{(I_2 + \Delta^2 J_2 + I_1) - I_0},$$

where the various integrals are as follows:

$$\tilde{I}_1(q, i\nu_m \rightarrow \omega + i0^+) = \int \frac{d^3k}{(2\pi)^3} I_1(k, q),$$

$$J_1(q, i\nu_m \rightarrow \tilde{\omega} + i0^+) = \int \frac{d^3k}{(2\pi)^3} I_2(k, q),$$

$$J_2(q, i\nu_m \rightarrow \omega + i0^+) = \int \frac{d^3k}{(2\pi)^3} I_2(k, q),$$

$$\tilde{I}_2(q, i\nu_m \rightarrow \omega + i0^+) = \int \frac{d^3k}{(2\pi)^3} (\tilde{E}^2 - x^2) I_2(k, q).$$

The method of calculation of various integrals is described in appendix A. Eliminating $\Gamma^{(1)}$ from equations (31) and (32), we have

$$\Gamma^{(2)} = i\tilde{\Delta}\tilde{\omega} \left[ 1 - \frac{(\tilde{\omega} J_1)^2}{D_+D_-} \right]^{-1} \left[ \frac{J_2}{D_+} + \frac{2J_1^2}{D_+D_-} \right] \Gamma^{(3)} ,$$

where $D_\pm = (\tilde{I}_1 + \tilde{I}_2 - I_0) \pm \tilde{\Delta}^2 J_2$. Here, we note that, although the integrals $\tilde{I}_1, \tilde{I}_2$ and $I_0$ have logarithmic divergence, this divergence does not pose any problem. Because, at the end, all the divergences are exactly cancelled out and thus all the vertex terms $\Gamma^{(i)}$’s and the response function remain finite. The integrals $J_1$ and $J_2$ are finite. From equation (26), we have

$$\Gamma^{(0)} = \gamma_0 \frac{1}{1 + \kappa B},$$

where

$$B = \tilde{I}_2 - \tilde{\Delta}^2 J_2 - \tilde{I}_1. $$

On substitution of equations (31) and (32), we have

$$\Gamma^{(3)} = \gamma_3 \frac{1}{1 + \kappa F},$$

where

$$F = A + \left[ (\tilde{\Delta}\tilde{\omega})^2 J_2 - \frac{2(\tilde{\Delta}\tilde{\omega} J_1)^2}{D_-} \right] \left[ \frac{J_2}{D_+} + \frac{2J_1^2}{D_+D_-} \right] \left[ 1 - \frac{(\tilde{\omega} J_1)^2}{D_+D_-} \right]^{-1} + \frac{(2\tilde{\Delta} J_1)^2}{D_-} $$

with

$$A = \tilde{I}_2 - \tilde{\Delta}^2 J_2 - \tilde{I}_1$$

After having carried out the Matsubara frequency sum and analytic continuation, we can write the response function as (see appendix B)

$$\chi(q, \omega) = 2B\gamma_0 \Gamma^{(0)} + 2A\gamma_3 \Gamma^{(3)} - 2i\tilde{\omega} \tilde{\Delta} J_2 \gamma_3 \Gamma^{(2)} + 4\tilde{\Delta} J_1 \gamma_3 \Gamma^{(1)}.$$
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Figure 1. Gap $\Delta$ (in units of $\epsilon_F$) is plotted as a function of temperature $k_B T$ (in units of $\epsilon_F$) for $k_F |a_s| = 2$ (a) and $k_F |a_s| = 4$ (c), while plots (b) and (d) exhibit the variation of the chemical potential $\mu$ as a function of temperature for $k_F |a_s| = 2$ and $k_F |a_s| = 4$, respectively.

On replacing $\Gamma^{(2)}$ and $\Gamma^{(1)}$ in terms of $\Gamma^{(3)}$, we have

$$\chi(q, \omega) = 2B\gamma_0 \Gamma^{(0)} + 2F\gamma_3 \Gamma^{(3)}$$

$$= \frac{2B}{1 + \kappa_s B \gamma_0^2} + \frac{2F}{1 + \kappa_s F \gamma_3^2}$$

(44)

Now the dynamic structure factor can easily be written as

$$S(q, \omega) = -\frac{1}{\pi} \left[ 1 + n_B(\omega) \right] \text{Im}[\chi(q, \omega)].$$

$$= -\frac{2}{\pi} \left[ 1 + n_B(\omega) \right] \left[ \gamma_0^2 \text{Im}(B) \left| \frac{1}{1 + \kappa_s B} \right|^2 + \gamma_3^2 \text{Im}(F) \left| \frac{1}{1 + \kappa_s F} \right|^2 \right].$$

(45)

4. Results and discussions

Before we elaborate our results, we note a few pertinent points. First, in calculating the response function, unlike weak-coupling case, all the energy integrations are carried out over the entire energy range, i.e. $\xi$ ranges from $-\mu$ to $\infty$. Second, as regards the vertex correction in light scattering, our formalism of calculating the response function can account for any arbitrary polarization of the incident light. Third, we have devised a procedure whereby we carry out the angular integrations by parts considering $\omega$ as a complex parameter $z$ as described in appendix A. This leads to trigonometric functions of $z$, which are then easily analytically continued using the limit $z \to \omega + i0^+$. This is unlike the usual approach where the functions such as $1/|\omega \pm (E_{\pm}) + i0^+|$ are first separated into real and imaginary parts. The imaginary part is a delta function, the energy and angular integrations over which require finding the roots of a complicated polynomial equation.

For all our numerical illustrations, the coupling strength $\kappa_s$ is taken to be negative since we are interested in Cooper-pairing regime only. Figure 1 shows the temperature variation of
Figure 2. (a): Dynamic structure factor $S(\omega, q)$ (in dimensionless units) is plotted as a function of energy transfer $\omega$ (in units of $\epsilon F$) for $q = 0.3 k_F$, $k_F |a_i| = 2.0$, $k_B T = 0.009 \epsilon_F$ ($T_c = 0.0265 \epsilon_F$) for unpolarized light with $\gamma_{11} = \gamma_{22} = \gamma$, that is, $\gamma_0 = 0$ and $\gamma_3 = \gamma$. (b) Same as in (a), but for circularly polarized light with $\gamma_{11} = 0$ and $\gamma_{22} \neq 0$, that is, $|\gamma_0| = |\gamma_1| = \gamma_{22}/2$. Plots (c) and (d) are the counterparts of (a) and (b), respectively; but for $k_B T = 0.021$.

Let us next discuss the regime of collective excitation of density–density fluctuations for energies $\omega < 2 \Delta$. This regime is characterized by the long-wave mode of vibration known as Bogoliubov–Anderson (BA) phonon which results from the vertex correction. The general expression for $\chi$ as given by equation (45) has two parts which are proportional to $\gamma_0^2$ and $\gamma_3^2$, respectively. Since the real part of $B$ is always negative, the first part is always finite. When $\Gamma^{(0)} \rightarrow \gamma_0$, $\Gamma^{(3)} \rightarrow \gamma_3$ and $\Gamma^{(1)} \rightarrow 0$ with $i = 1, 2$, that is, when the vertex correction is neglected, the response function reduces to

$$\chi(q, \omega) = 2B\gamma_0^2 + 2A\gamma_3^2$$

which has no poles. However, when the vertex correction is added, there arises the pole of $\chi$ which is given by

$$F = \frac{1}{|\kappa_i|}.$$ 

In the limit $|a_i| \rightarrow 0$, the pole is determined by $D_+ = 0$. The real part of the pole will correspond to the phonon energy while the negative imaginary part corresponds to the damping of the mode. Only in the low $q$-regime and $T < T_c$, the BA mode will be well defined. For larger temperatures (but $< T_c$), this mode will be ill defined because of the large Landau damping which occurs due to its coupling with thermally excited quasiparticles. At zero
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Figure 3. $S(\omega, q)$ as a function of $\omega$ for two different values of $q$. The other parameters are $k_F |a_1| = 4.0$, $k_B T = 0.055 \epsilon_F$ ($T_F = 0.124 \epsilon_F$, $\Delta = 0.203 \epsilon_F$), $\gamma_{11} = \gamma_{22} = 1$.  

$$S(\omega, q)$$ as a function of $\omega/\epsilon_F$ for two different values of $q = 0.3k_F$ and $q = 0.5k_F$. $x_{10} = 4$, $k_B T = 0.055 \epsilon_F$ ($T_F = 0.124 \epsilon_F$, $\Delta = 0.203 \epsilon_F$), $\gamma_{11} = \gamma_{22} = 1$. 

Figure 4. Quasiparticle excitation spectrum ($\omega > 2\Delta$) for a uniform Fermi superfluid at two different temperatures $k_B T = 0.115 \epsilon_F$ (a, b) and $k_B T = 0.075 \epsilon_F$ (c, d), but for the same $k_F |a_1| = 4.0$. The plots (a) and (c) are for unpolarized light with $\gamma_{11} = \gamma_{22} = 1$, while plots (b) and (d) are for circularly polarized light with $\gamma_{11} = 0$. The dashed curve refers to $q = 0.4k_F$ and the solid one to $q = 0.6k_F$ in all the plots. $\Delta$ is 0.098 $\epsilon_F$ and 0.188 $\epsilon_F$ at temperatures 0.115 $\epsilon_F$ and 0.075 $\epsilon_F$, respectively. 

Temperature, in the limit $a_s \to 0^-$, we have reproduced the standard result 

$$\omega_{BA} = \frac{1}{\sqrt{3}} \hbar q v_F$$  (48) 

as calculated in appendix C. The BA phonon appears as a resonance in the spectrum of dynamic structure factor as is demonstrated in figures 2 and 3. Comparing figure 2(a) with figure 2(c), we note that as the temperature is increased, the spectrum becomes broadened. The resonance energy does not depend on the state of polarization of light. However, the peak and width of the resonance may depend on the polarization of light. Figure 4 shows the quasiparticle excitation
spectrum when $\omega$ exceeds the pair-breaking energy $2\Delta$. As temperature is decreased, the peak makes a positive shift because of an increase in $\Delta$. For larger $q$ values, the peak is larger. At lower temperatures, the polarization of light significantly affects the spectrum as can be noted by comparing figure 4(c) with 4(d).

5. Conclusion

In conclusion, we have derived the finite temperature response function of the superfluid Fermi gas due to scattering of polarized light. The response function takes into account vertex correction due to final-state interactions. We have presented selective results on dynamic structure factor (DSF) deduced from the response function with the aid of generalized fluctuation–dissipation theorem. The in-gap collective mode of the density fluctuation, known as Bogoliubov–Anderson phonon, appears as a strong resonance in the plot of DSF as a function of energy transfer $\omega$ at low momentum transfer $q$. As the temperature increases, the width of the resonance increases due to Landau damping. At large scattering lengths ($k_F|a_s| > 1$), resonance may occur at a finite and appreciable value of $q$. Polarization-selective light scattering may be useful in exciting the BA mode. We have also presented results on the single-particle excitation spectrum when Cooper-pairs are broken ($\omega > 2\Delta$). As the temperature is decreased, the peak of the spectrum makes a positive shift due to increase of pair-breaking energy ($2\Delta$). We have also shown the effect of light polarization on collective as well as single-particle excitations. Particularly important is the positive shift of the peak of the single-particle excitation spectrum at lower temperatures due to circularly polarized light as compared to that due to unpolarized light (figure 4). Furthermore, the higher the momentum transfer, the sharper is the single-particle excitation spectrum. As pointed out earlier [27], in stimulated light scattering, the momentum transfer may be large and is comparable with the Fermi momentum. Since in the polarization-selective stimulated light scattering, only the atoms with the same spin state are predominantly scattered [28], this kind of light scattering with large momentum transfer may be useful in detecting the gap energy and its temperature dependence. In the present paper, our study is confined to uniform FS only. Local density approximation (LDA) may be applied for studying single-particle excitation of a trapped FS at a large momentum transfer as shown in [28]. However, for studying the collective mode of a trapped system, LDA may be a bad approximation. We hope to address this problem in our future communication.

Appendix A

Here, we first calculate the four $T_{ij}$ terms as expressed in equations (21)–(24) and then describe the method of calculation of various integrals. We express these as follows:

$$T_{11} = \beta^{-1} \sum_n \frac{1}{i(\omega_n + v_m/2) + E_{k-}} \times \frac{1}{i(\omega_n - v_m/2) + E_{k'}}$$

$$= \beta^{-1} \frac{1}{E_{k-} - E_{k'} + iv_m} \sum_n \left[ \frac{1}{i(\omega_n - v_m/2) + E_{k'}} - \frac{1}{i(\omega_n + v_m/2) + E_{k'}} \right]$$

$$= \frac{n_F(-E_{k-}) - n_F(-E_{k'})}{E_{k-} - E_{k'} + iv_m}. \quad (A.1)$$
Similarly,

\[ T_{12} = \frac{n_F(E_{k'}) - n_F(-E_{k'})}{E_{k'} + E_{k'} + iv_m} \]  \hspace{1cm} (A.2)  
\[ T_{21} = \frac{n_F(-E_{k'}) - n_F(E_{k'})}{-E_{k'} - E_{k'} + iv_m} = \frac{n_F(E_{k'}) - n_F(-E_{k'})}{E_{k'} + E_{k'} - iv_m} \]  \hspace{1cm} (A.3)  
\[ T_{22} = \frac{n_F(E_{k'}) - n_F(E_{k'})}{E_{k'} - E_{k'} - iv_m}. \]  \hspace{1cm} (A.4)  

One can write \( n_F(E) = 1/2 - \tanh(\beta E/2)/2 \) and \( n_F(-E) = 1/2 + \tanh(\beta E/2)/2 \). Using these relations, we obtain equations (21)–(24).

Let us now consider the integrals \( I_i \) and \( J_i \). Setting \( z = iv_m \), we can write

\[ I_1(q, z) = \frac{1}{2(\pi)^2} \int \int \sin \theta \, d\theta \, k^2 \, dk \, I_1(k, q, z) \]
\[ = \frac{1}{4(2\pi)^2} \int k^2 \, dk \int_{-1}^{1} dy \, R_+ (k, q, y, z) \]  \hspace{1cm} (A.5)  
\[ J_2(q, z) = \frac{1}{4(2\pi)^2} \int k^2 \, dk \int_{-1}^{1} dy \, \frac{1}{E_z E_1} \, R_- (k, q, y, z), \]  \hspace{1cm} (A.6)  

where

\[ R_\pm (k, q, y, z) = \frac{T_\pm (k, q)(E_2 + E_1)}{(E_2 + E_1)^2 - z^2} \pm \frac{T_\mp (k, q)(E_2 - E_1)}{(E_2 - E_1)^2 - z^2}. \]  \hspace{1cm} (A.7)  

Here

\[ T_\pm = \tanh(\beta E_2 / 2) \pm \tanh(\beta E_1 / 2). \]  \hspace{1cm} (A.8)  

Here \( y = \cos \theta \). Similarly, we can express

\[ J_1(q, z) = \frac{1}{4(2\pi)^2} \int k^2 \, dk \int_{-1}^{1} dy \, \frac{\mathcal{E}}{E_z E_1} \, R_- (k, q, y, z) \]  \hspace{1cm} (A.9)  
\[ I_2(q, z) = \frac{1}{4(2\pi)^2} \int k^2 \, dk \int_{-1}^{1} dy \, \frac{\mathcal{E}^2 - x^2}{E_z E_1} \, R_- (k, q, y, z). \]  \hspace{1cm} (A.10)  

In the limit \( T \to 0 \), we have \( T_+ = 2 \) and \( T_- = 0 \). Now,

\[
\bar{I}_1 = \frac{1}{4(2\pi)^2} \int k^2 \, dk \int_{-1}^{1} dy \\
\times \left[ T_+(E_1 + E_2)((E_2 - E_1)^2 - z^2) + T_-(E_2 - E_1)((E_1 + E_2)^2 - z^2) \right] \\
= \frac{1}{4(2\pi)^2} \int k^2 \, dk \int_{-1}^{1} dy \left( E_2^2 - E_1^2 \right)^2 - 2z^2 \left( E_1^2 + E_2^2 \right) + z^4 \\
\times \left[ E_2 \tanh \left( \frac{\beta E_2}{2} \right) \left( E_2^2 - E_1^2 - z^2 \right) - E_1 \tanh \left( \frac{\beta E_1}{2} \right) \left( E_2^2 - E_1^2 + z^2 \right) \right].
\]
Substituting $E_2^2 - E_1^2 = 4\mathcal{E}x$ and $E_1^2 + E_2^2 = 2\mathcal{E}^2 + 2x^2 + 2\Delta^2$, the above equation can then be expressed as

\[
\tilde{I}_1 = \frac{3\rho_0}{8\epsilon_0 \tilde{q}} \int d\mathcal{E} \int_{-\tilde{q}}^{\tilde{q}} dx \frac{1}{z^2(z^2 - 4\mathcal{E}^2 + 4\Delta^2) + 4(4\mathcal{E}^2 - z^2)x^2}[\frac{-\Delta^2}{2\mathcal{E}} - 4\mathcal{E}x\mathcal{F}_-]
\]

\[
= \frac{3\rho_0}{16\epsilon_0 \tilde{q}} \int d\mathcal{E} \int_{-\tilde{q}}^{\tilde{q}} dx \frac{1}{z^2(\mathcal{E}^2 - Z_0^2)} \int_0^{\tilde{q}} d\mathcal{K} \Phi(x) \frac{1}{1 + \nu^2 x^2},
\]

(A.11)

where

\[
\nu^2 = \frac{z^2 - 4\mathcal{E}^2}{z^2(\mathcal{E}^2 - Z_0^2)}
\]

(A.12)

\[
Z_0^2 = \frac{1}{4}(z^2 - 4\Delta^2)
\]

(A.13)

\[
\Phi(x) = z^2 \mathcal{F}_x + 4\mathcal{E}x \mathcal{F}_-
\]

(A.14)

with $\mathcal{F}_x = E_1 \tanh(\beta E_1/2) \pm E_2 \tanh(\beta E_2/2)$. Here $\rho_0 = k^3/4\pi^2$ is the number density of noninteracting single-spin Fermi gas. We further note that $E_1(x) = E_2(-x)$ implying $\mathcal{F}_x(x) = \pm \mathcal{F}_x(-x)$. Let us carry out $x$-integration by parts, assuming $\Phi(x)$ and $(1 - \nu^2 x^2)^{-1}$ be the first and second integrand, respectively. Then, we have

\[
I_0 = \int_0^{\tilde{q}} d\mathcal{K} \Phi(x) \frac{1}{1 + \nu^2 x^2}
\]

\[
= \frac{1}{\nu} \left[ \Phi(\tilde{q} \mathcal{K}) \tan^{-1}((\nu \tilde{q} \mathcal{K})) - \int_0^{\tilde{q}} d\mathcal{K} \Phi'(x) \tan^{-1}(\nu x) \right].
\]

(A.15)

Let us now consider the analytic continuation by taking the limit $z \to \tilde{\omega} + i0^+$, where $\tilde{\omega}$ is the energy transfer. There are two regimes of excitations: (i) Quasiparticle regime: $\tilde{\omega} > 2\Delta$ and (ii) collective oscillation regime: $\tilde{\omega} < 2\Delta$. We first concentrate on the regime of quasiparticle excitation, that is, $\tilde{\omega} > 2\Delta$. In this regime, $\mathcal{E}_0 > 0$. Let $\phi = \tan^{-1}(\nu x)$ and $\mathcal{E}_0 = \sqrt{\tilde{\omega}^2 - 4\Delta^2}/2$. We have three cases for consideration of analytic continuation:

1. For $|\mathcal{E}| \leq \mathcal{E}_0$, we have

\[
\nu \to i \nu_1 = \frac{1}{\tilde{\omega}} \sqrt{\tilde{\omega}^2 - 4\mathcal{E}^2} \quad \mathcal{E}_0^2 - \mathcal{E}^2
\]

(A.16)

\[
\phi \to i \tanh^{-1}(\nu_1 x), \quad \nu_1 x \leq 1
\]

(A.17)

\[
\phi \to \frac{\pi}{2} + i \tanh^{-1} \left( \frac{1}{\nu_1 x} \right), \quad \nu_1 x > 1.
\]

(A.18)

2. For $\mathcal{E}_0 < |\mathcal{E}| \leq \tilde{\omega}/2$, we have

\[
\nu \to \nu_2 = \frac{1}{\tilde{\omega}} \sqrt{\tilde{\omega}^2 - 4\mathcal{E}^2} \quad \mathcal{E}_0^2 - \mathcal{E}^2
\]

(A.19)

\[
\tan \phi \to \nu_2 x.
\]

(A.20)
For $|\mathcal{E}| > \tilde{\omega}/2$, we have

\[
\nu \rightarrow iv_3 = \frac{1}{\tilde{\omega}} \sqrt{\frac{4\mathcal{E}^2 - \tilde{\omega}^2}{\mathcal{E}^2 - \mathcal{E}_0^2}} \tag{A.21}
\]

\[
\phi \rightarrow i \tanh^{-1}(v_3x), \quad v_3x \leq 1 \tag{A.22}
\]

\[
\phi \rightarrow \frac{\pi}{2} + i \tanh^{-1}\left(\frac{1}{v_3x}\right), \quad v_3x > 1, \tag{A.23}
\]

After having carried out the angular integration and analytic continuation, the energy integration is carried out numerically. Let us next consider the integral $J_2$. Towards this end, we have

\[
\frac{\mathcal{R}_-}{E_1 E_2} = \frac{2}{\pi^2} \frac{\nu_3 x}{(E_2 - E_0^2)(1 + v^2 x^2)}, \tag{A.24}\]

where $\mathcal{J}_\pm(x) = \tanh[\beta E_1/2]/E_2 \pm \tanh[\beta E_1/2]/E_1$. Using this, we can write

\[
J_2 = \frac{3\nu_0}{16\pi f \tilde{q}} \int d\mathcal{E} \frac{1}{z^2(E^2 - \mathcal{E}_0^2)} \int_0^{\tilde{q}k} dx \frac{\Psi(x)}{\nu_3 x}, \tag{A.25}\]

where

\[
\Psi(x) = z^2 \mathcal{J}_+ + 4\mathcal{E}_x \mathcal{J}_-. \tag{A.26}\]

Thus the integral $J_2$ can be evaluated exactly the same way as in the case of $\tilde{I}_1$. Similarly, we can find the other two integrals $\tilde{I}_2$ and $J_1$.

**Appendix B**

We can rewrite the response equation

\[
\chi(q, \omega) = \int \frac{d^3k}{(2\pi)^3} \delta^{-1} \sum_n \frac{1}{[(\omega_n + v_m)^2 + E_2^2][\omega_n^2 + E_1^2]} \times \text{Tr}[\gamma K [i(\omega_n - v_m/2)\tau_0 + \xi_k \tau_3 + \Delta \tau_1]] \tag{B.1}\]

where

\[
K = i((\omega_n + v_m/2)) \sum_{j=0}^3 \Gamma^{(j)} \tau_j + \xi_k \left\{ \Gamma^{(0)} \tau_3 + \Gamma^{(3)} \tau_0 + i \sum_{j=1,2} \epsilon_{ij} \Gamma^{(j)} \tau_{3-j} \right\} + \Delta \left\{ \Gamma^{(0)} \tau_1 + \Gamma^{(1)} \tau_0 + i \sum_{j=2,3} \Gamma^{(j)} \epsilon_{1j} \tau_{n_j} \right\}, \tag{B.2}\]

where $n_j = |1 + \epsilon_{ij}|$ and $\epsilon_{ij} = -\epsilon_{ji} = 1$ if $(i, j)$ is $(1, 2)$, or $(2, 3)$, or $(3, 1)$. Now,

\[
K \times [i(\omega_n - v_m/2)\tau_0 + \xi_k \tau_3 + \Delta \tau_1] = i(\omega_n - v_m/2) \mathcal{K} + \xi_k \cdot i((\omega_n + v_m/2)) \left\{ \Gamma^{(0)} \tau_3 + \Gamma^{(3)} \tau_0 + i \sum_{j=1,2} \epsilon_{j3} \Gamma^{(j)} \tau_{3-j} \right\} + \xi_k \xi_k \left\{ \Gamma^{(0)} \tau_0 + \Gamma^{(3)} \tau_3 - \sum_{j=1,2} \Gamma^{(j)} \tau_j \right\}.
\]
Here we turn our attention to the collective oscillation regime where $E_0^2 < 0$. We do not make any attempt to evaluate analytically the collective mode for any interaction strength $(k_F|C_i|)$ at the finite temperature. However, it is possible to calculate analytically the BA mode energy in the weak-coupling limit $k_F|a_i| 	o 0$ at zero temperature. Towards this end, let us first calculate $D_{r} (z 	o \omega + i0^{+})$. Restricting $\chi < 1$, we expand the functions $E_{\pm} = E_{2,1}(x)$ and $1/E_{\pm}(x)$ up to second order in $x$:

\begin{equation}
E_{\pm} \simeq E + \frac{x^2 \pm 2Ex}{2E} - \frac{\tau E^2 x^2}{2E^3}.
\end{equation}

\begin{equation}
\frac{1}{E_{\pm}} \simeq \frac{1}{E} \left[ 1 - \frac{x^2 \pm 2Ex}{2E^2} + \frac{3\tau E^2 x^2}{2E^4} \right].
\end{equation}

At zero temperature, we obtain

\begin{equation}
\Phi(x \simeq 0) \simeq \phi^2 \left( 2E + \frac{x^2}{E} \right) - 8\phi^2 \frac{x^2}{E} - \frac{\phi^2 \tau E^2 x^2}{E^3},
\end{equation}

\begin{equation}
\Psi(x \simeq 0) \simeq \frac{1}{E} \left[ \phi^2 \left( 2 - \frac{x}{E^2} \right) - 8\phi^2 \frac{x^2}{E^2} + \frac{3\phi^2 \tau E^2 x^2}{E^4} \right].
\end{equation}

We can write $I_1 + I_2 + \Delta^2 J_2 \simeq I_\chi + I_\sigma$, where

\begin{equation}
I_\chi = \frac{3\rho_0}{4\pi} \int_{|\vec{k}| \lesssim \omega/2} d\vec{E} \left[ \frac{E}{(E^2 - \vec{E}_0^2)} \tan^{-1}(v_2 \vec{k}) \right.
\end{equation}
Here $I_\omega$ is real, but $I_\nu$ has both real and imaginary parts. We now calculate

$$\text{Re}[I_\omega] = \frac{3\rho_0}{4\epsilon_F} \int_{\epsilon_\nu - |\epsilon| - \bar{\omega}/2} d\epsilon \left[ \frac{E}{(\epsilon^2 - \epsilon_\nu^2)} \tanh^{-1}(\nu \bar{q} \bar{k}) \right]$$

$$= \frac{1}{\bar{\omega}^2(\epsilon^2 - \epsilon_\nu^2)} \left( \frac{\bar{\omega}^2 + 8\epsilon^2}{2\epsilon} - 2\frac{\bar{\omega}^2\epsilon^2}{E^3} \right) \left( -k \bar{v}_\perp^2 + \frac{\tanh^{-1}(\nu \bar{q} \bar{k})}{\bar{v}_\perp} \right)$$

$$+ \frac{3\rho_0}{4\epsilon_F} \int_{|\epsilon| > \epsilon_\nu} d\epsilon \left[ \frac{E}{(\epsilon^2 - \epsilon_\nu^2)} \tanh^{-1}[1/(\nu \bar{q} \bar{k})] \right]$$

$$= \frac{\bar{\omega}^2 + 8\epsilon^2}{2E\bar{\omega}^2(\epsilon^2 - \epsilon_\nu^2)} \left( -\frac{k}{\bar{v}_\perp^2} + \frac{\tanh^{-1}[1/(\nu \bar{q} \bar{k})]}{\bar{v}_\perp} \right)$$

$$= \frac{3\rho_0 \bar{k}_\mu}{4\epsilon_F} \left[ \left( \frac{1}{2} - \frac{\epsilon^2}{3\bar{\omega}^2} \right) \left( \frac{\bar{\omega}}{\sqrt{1 - \bar{\omega}^2}} \right) (\pi - 2 \sin^{-1}(\sqrt{1 - \bar{\omega}^2})) \right]$$

$$= \frac{(\bar{q} \bar{k}_\mu)^2}{3\bar{\omega}^3} \left[ \frac{2\bar{\omega} \Delta \sqrt{1 - \bar{\omega}^2} - 4\Delta^2(1 - 2\bar{\omega}^2) \sin^{-1}(\bar{\omega})}{(1 - \bar{\omega}^2)^{3/2}} \right]$$

$$+ \frac{(\bar{q} \bar{k}_\mu)^2}{3} \left( \frac{4\Delta}{\bar{\omega}^3} (\bar{\omega} - \sqrt{1 - \bar{\omega}^2} \sin^{-1}(\bar{\omega})) \right),$$

where $\bar{\omega} = \bar{\omega}/(2\Delta)$. For $\bar{\omega} \ll 1$, we can expand

$$\sin^{-1}(\sqrt{1 - \bar{\omega}^2}) \approx \sin^{-1}(1 - \frac{1}{2}\bar{\omega}^2) = \frac{\pi}{2} - \bar{\omega} - \frac{1}{12}\bar{\omega}^3 + \cdots.$$  

Similarly, expanding $\sin^{-1}(\bar{\omega})$ and retaining the terms of lower order in $\bar{\omega}$, we can approximate the square bracketed part of equation (C.8) as

$$\frac{1}{\sqrt{1 - \bar{\omega}^2}} \left[ \bar{\omega}^2 + \frac{1}{12}\bar{\omega}^3 - \frac{(\bar{q} \bar{k}_\mu)^2}{3(2\Delta)^2} \left( 2 + \frac{1}{6}\bar{\omega}^2 \right) \right] \approx -\frac{8(\bar{q} \bar{k}_\mu)^2}{9(2\Delta)^2(1 - \bar{\omega}^2)^{3/2}} + \frac{2(\bar{q} \bar{k}_\mu)^2}{9(2\Delta)^2}.$$  

Furthermore, keeping the terms up to the second order in $\bar{\omega}$ and neglecting the product terms like $(\bar{q} \bar{k}_\mu)^2 \bar{\omega}^2$, the above expression reduces to

$$\bar{\omega}^2 = \frac{4}{3} \left( \frac{\bar{q} \bar{k}_\mu}{2\Delta} \right)^2.$$
Equating this to zero, we find the root \( \tilde{\omega} = \frac{2}{\sqrt{3}} \tilde{q} \tilde{k}_\mu \) which implies

\[
\omega = \frac{1}{\sqrt{3}} p_q v_\mu,
\]

where \( p_q = \hbar q \) and \( v_\mu = \hbar k_\mu / m \).

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