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Citation: *Physics of Plasmas* **24**, 074501 (2017); doi: 10.1063/1.4989375

View online: <http://dx.doi.org/10.1063/1.4989375>

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On the Raman instability in degenerate relativistic plasmas

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(Received 8 June 2017; accepted 21 June 2017; published online 10 July 2017)

The stimulated Raman scattering instability in a fully degenerate electron plasma is studied applying relativistic hydrodynamic and Maxwell equations. We demonstrated that the instability develops for weakly and strongly relativistic degenerate plasmas. It is shown that in the field of strong radiation, a degenerate relativistic plasma effectively responds as in the case of weak degeneracy. *Published by AIP Publishing.* [<http://dx.doi.org/10.1063/1.4989375>]

Astrophysical objects contain very intense radiation sources with their spectral content spanning from radio to γ -ray emission.¹ Most types of stars, interstellar gas, active galactic nuclei, distant quasars, *etc.*, are either comprised by a significant fraction of plasma or immersed in a plasma environment.² The nature of radiation emitted by such plasmas is usually interpreted in terms of bremsstrahlung and synchrotron mechanisms. However, under certain conditions, the collective parametric effects such as stimulated Raman scattering (SRS) instability could leave a definite signature on the radiation spectra.³ Within the framework of classical plasmas, the SRS has been studied in a variety of systems such as the laser-driven inertial confinement fusion, particle acceleration, and plasma heating and for probing the plasma parameters in laboratory conditions.^{4–10} Linear and nonlinear stages of SRS instability have been investigated for cold and relativistic hot plasmas for electromagnetic (EM) radiation of arbitrary intensity. A possible role of SRS instability in an astrophysical plasma is discussed in Refs. 11 and 12 and references therein; it is shown that induced scattering could significantly affect radiation from sources with high brightness temperatures.

For the extremely high density plasmas pertinent to astrophysical objects such as in the interior of white dwarfs, the neutron or the pre-supernova stars, and presumably, at the source of gamma ray bursts,² the classical approximation breaks down; the relevant plasma is degenerate requiring Fermi-Dirac statistics for a correct description. The Fermi energy of the degenerate electron gas is greater than the binding energy with the atomic nucleus, and as a consequence, all atoms are in an ionized state. The plasma number density is believed to be in the range from $N_0 = 10^{26} \text{cm}^{-3}$ to $N_0 = 10^{34} \text{cm}^{-3}$, and the matter behaves as a weakly coupled degenerate plasma provided that averaged interparticle distance is smaller than the thermal de Broglie wavelength.¹³ At such densities, the degenerate electron gas must be treated relativistically even when its “temperature” is nonrelativistic, in fact, even zero.

The Fermi energy associated with typical particles in a degenerate electron gas, $\epsilon_F = m_e c^2 (\gamma_F - 1)$ (the Fermi $\gamma_F = \sqrt{1 + p_F^2/m_e^2 c^2}$, $p_F = m_e c (n_R/n_c)^{1/3}$ ¹⁴), exceeds the rest mass energy for densities n_R greater than the critical density $n_c = m_e^3 c^3 / 3\pi^2 \hbar^3 = 5.9 \times 10^{29} \text{cm}^{-3}$.

In studies of nonlinear self-interactions of high frequency EM and plasma waves in relativistic degenerate electron (as well as electron-positron) plasmas,^{15–18} it has been shown such plasmas can support stable localized EM structures for an arbitrary level of degeneracy. The dynamics of these nonlinear EM structures could provide a theoretical basis for establishing the nature of the observed radiation. To the best of our knowledge, the SRS instability of EM radiation in degenerate relativistic plasmas has not been addressed so far. The SRS instability is a nonlinear parametric process when a powerful electromagnetic (EM) wave decays into a plasma wave and an EM wave. In an underdense plasma with $\omega_0 > 2\omega_{e0}$, where ω_0 is the mean frequency of carrier EM pulses and $\omega_{e0} = (4\pi e^2 N_0/m_e)^{1/2}$ is the electron plasma frequency, the SRS results in two EM sidebands upshifted and downshifted by the plasma frequency. For a relativistic degenerate electron gas, the radiation frequency at which SRS could be relevant falls into soft or even hard X-ray bands. Consequently, it could leave definite footprints in the spectra of astrophysical X-ray sources.

Our study of the SRS instability is based on the Maxwell equations and a relativistic electron plasma fluid model. We will work in terms of the familiar vectorial form of the fully covariant equations (see, for example, Refs. 19–21) that translate as (in the Coulomb gauge $\nabla \cdot \mathbf{A} = 0$)

$$\frac{\partial^2 \mathbf{A}}{\partial t^2} - c^2 \Delta \mathbf{A} + c \frac{\partial}{\partial t} (\nabla \varphi) - 4\pi c \mathbf{J} = 0, \quad (1)$$

$$\Delta \varphi = 4\pi e (n_R \gamma - N_0), \quad (2)$$

where the vector potential \mathbf{A} and the scalar potential φ are the appropriate components of the EM four vector $A^\mu = [\varphi, \mathbf{A}]$ and $\mathbf{J} = -en_R \mathbf{p}/m_e$, and $-ec\gamma n_R$ constitutes the electron four current J^μ .

The ions just provide a neutralizing background and have a density N_0 in their rest frame which is also taken to be the fiduciary/laboratory frame. The “rest” frame electron density, n_R , is a Lorentz scalar and is related to the laboratory frame density N by the relation $N = \gamma n_R$. The relativistic $\gamma = \sqrt{1 + \mathbf{p}^2/m_e^2 c^2}$ is determined from the vector part \mathbf{p} of the four momentums.

In the plasma with zero generalized vorticity $\boldsymbol{\Omega} = \nabla \times (G\mathbf{p} - e\mathbf{A}/c) = 0$, the fluid equations for the electrons can be written as (see Ref. 15 for details) follows:

$$\frac{\partial}{\partial t} \left(G\mathbf{p} - \frac{e}{c}\mathbf{A} \right) + \nabla (m_e c^2 G\gamma - e\varphi) = 0, \quad (3)$$

$$\frac{\partial}{\partial t} n_R \gamma + \nabla \cdot (n_R \mathbf{p} / m_e) = 0. \quad (4)$$

In Eq. (3), the “thermal” mass enhancement factor G is defined as $G = w / (n_R m_e c^2)$, where w is the enthalpy per unit volume. The effective mass factor is, generally, a non-trivial function of plasma temperature and density for both classical Maxwell–Jüttner and the quantum Fermi–Jüttner statistics. However, if the thermal energy of the plasma electrons is much lower than the Fermi energy, the plasma may be treated as cold (i.e., having zero temperature). The zero temperature approximation will be adequate even for temperatures of order 10^9 K, allowing a particularly simple expression

$$G = \gamma_F = \sqrt{1 + (n_R / n_c)^{2/3}}.$$

For an EM wave propagating in the z direction (all dynamic variables vary only in z and time t), the vector potential has just transverse components $\mathbf{A} = (\mathbf{A}_\perp, 0)$. For this simplified 1-D propagation, Eq. (3) can be readily integrated to yield

$$\mathbf{p}_\perp = \frac{\mathbf{A}_\perp}{\gamma_F}, \quad (5)$$

where

$$\gamma_F = \sqrt{1 + (R_0 n)^{2/3}} \quad (6)$$

and

$$\gamma = \sqrt{1 + \frac{\mathbf{A}_\perp^2}{\gamma_F^2} + p_z^2} \quad (7)$$

in terms of the normalized variables $\mathbf{A}_\perp = (e\mathbf{A}_\perp / m_e c^2)$, $\varphi = (e\varphi / m_e c^2)$, $\mathbf{p} = (\mathbf{p} / mc)$, and $n = (n_R / N_0)$.

The parameter $R_0 = (N_0 / n_c)$ measures the strength of plasma degeneracy: for $R_0 \ll 1$, the plasma is nonrelativistic, while for $R_0 \geq 1$, the plasma is in the relativistic degenerate state. Note that the dimensionless Fermi momentum now reads as $p_F = (R_0 n)^{1/3}$.

The Maxwell equation (1) is reduced to the wave equation

$$\frac{\partial^2 \mathbf{A}_\perp}{\partial t^2} - c^2 \frac{\partial^2 \mathbf{A}_\perp}{\partial z^2} + \Omega_e^2 \mathbf{A}_\perp = 0, \quad (8)$$

where $\Omega_e = (4\pi e^2 n_R / m_e \gamma_F)^{1/2}$ is the frame independent plasma frequency and is a Lorentz scalar. In the standard literature, it is conventional to define the frequency $\omega_{e0} = (4\pi e^2 N_0 / m_e)^{1/2}$ corresponding to the density in the lab frame; the latter is related to the frame-invariant frequency through $\Omega_e = \omega_{e0} (n / \gamma_F)^{1/2}$, i.e., the invariant frequency, in this case, is seen as the lab frame frequency reduced by the relativistic Fermi effects. There are lots of subtle considerations in creating a strictly Lorentz invariant theory, and the reader is referred to Refs. 19–21 for a deeper discussion.

The system of equations describing the longitudinal motion of the electron plasma is given by

$$\frac{\partial}{\partial t} \gamma_F p_z + c \frac{\partial}{\partial z} (\gamma \gamma_F - \varphi) = 0, \quad (9)$$

$$\frac{\partial}{\partial t} \gamma n + c \frac{\partial}{\partial z} (n p_z) = 0, \quad (10)$$

$$\frac{\partial^2 \varphi}{\partial z^2} = \frac{\omega_{e0}^2}{c^2} (n \gamma - 1). \quad (11)$$

We next carry out a stability analysis for the circularly polarized EM waves. The monochromatic pump EM wave with frequency ω_0 and wave number k_0 is described as

$$\mathbf{A}_\perp = \frac{1}{2} (\hat{\mathbf{x}} + i\hat{\mathbf{y}}) A \exp(-i\omega_0 t + ik_0 z) + c.c. \quad (12)$$

Here, $\hat{\mathbf{x}}$ and $\hat{\mathbf{y}}$ are the unit vectors; $A = a \exp(i\psi)$, where a and ψ are the real valued amplitude and phase. The unperturbed state of the plasma is characterized by a purely transverse EM mode with constant amplitude $a = a_0$ ($p_z = 0 = \varphi$) and a constant co-moving density of electrons n_0 . This density is related to the lab frame density by the relation $n_0 = 1 / \gamma_0$ (in units $n_0 = N_0 / \gamma_0$), where $\gamma_0 = \sqrt{1 + a_0^2 / \gamma_{F0}^2}$ and $\gamma_{F0} = \sqrt{1 + (R_0 / \gamma_0)^{2/3}}$.

The dispersion relation that follows from Eqs. (8)–(12) reads

$$\omega_0^2 = k_0^2 c^2 + \Omega_{e0}^2, \quad (13)$$

where $\Omega_{e0} = \omega_{e0} (n_0 / \gamma_{F0})^{1/2} = \omega_{e0} / \gamma_m^{1/2}$ is the relativistically modified electron plasma frequency and $\gamma_m = \gamma_0 \gamma_{F0}$. The expression for γ_m can be written as $\gamma_m = \sqrt{1 + p_{F0}^2 + a_0^2}$, where $p_{F0} = (R_0 / \gamma_0)^{1/3}$ is the Fermi momentum. In the weakly degenerate case, when $R_0 \ll 1$ and $p_{F0} \ll a_0$, Eq. (13) coincides with the dispersion relation for cold classical plasmas with $\Omega_{e0} = \omega_{e0} / (1 + a_0^2)^{1/2}$, while for the arbitrary R_0 and in the absence of the pump wave ($a_0 = 0$), the modified plasma frequency reads $\Omega_{e0} = \omega_{e0} / (1 + R_0^{2/3})^{1/2}$.

Most spectacular manifestations of plasma degeneracy, however, occur for radiation with relativistically large amplitudes, i.e., when $a_0 \gg 1$, and the radiative modification of the effective electron-mass becomes comparable to or dominant over the degeneracy modification (see Ref. 20 for a detailed discussion on the radiative renormalization of the electron mass). For extreme relativistic amplitudes, the effective γ simplifies to $\gamma_m \simeq \sqrt{1 + a_0^2}$, essentially the expression for a nondegenerate plasma. For arbitrary values of a_0 and R_0 (measuring the degeneracy strength), the relativistic factor γ_0 can be found from the following implicit relation:

$$(\gamma_0^2 - 1) \left(1 + \left(\frac{R_0}{\gamma_0} \right)^{2/3} \right) = a_0^2 \quad (14)$$

that can be readily solved. In Fig. 1, we plot γ_0 vs a_0 for different values of R_0 . The dashed part of each curve corresponds to $p_{F0} < 1$. For larger values of R_0 , stronger fields are required to achieve large values of γ_0 .

A short message of the preceding results is that in the field of strong EM waves (see Fig. 1), the effects of electron degeneracy become subdominant and the plasma responds as a weakly degenerate system.

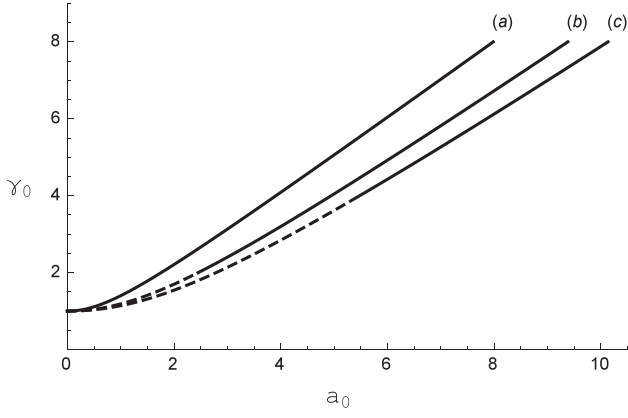


FIG. 1. Dependence of γ_0 on the EM field strength a_0 for different levels of degeneracy: (a) $R_0 = 0.01$, (b) $R_0 = 2$, and (c) $R_0 = 4$.

To investigate the stability of the ground wave solution, we introduce the small perturbations in the system $f = f_0 + \delta f$ ($f_0 \gg \delta f$), where $f = (a, \psi, n, p_z, \varphi)$. Neglecting higher order terms, Eqs. (8)–(11) can be reduced to the following system of coupled equations:

$$\left(2\omega_0 \frac{\partial}{\partial t} + 2k_0 c^2 \frac{\partial}{\partial z}\right) \delta\psi + \left(\frac{\partial^2}{\partial t^2} - c^2 \frac{\partial^2}{\partial z^2}\right) \frac{\delta a}{a_0} + \Omega_{e0}^2 \left(\frac{\delta n}{n_0} - \frac{\delta\gamma_F}{\gamma_{F0}}\right) = 0, \quad (15)$$

$$\left(2\omega_0 \frac{\partial}{\partial t} + 2k_0 c^2 \frac{\partial}{\partial z}\right) \frac{\delta a}{a_0} + \left(c^2 \frac{\partial^2}{\partial z^2} - \frac{\partial^2}{\partial t^2}\right) \delta\psi = 0, \quad (16)$$

$$\left[\frac{\partial^2}{\partial t^2} + \Omega_{e0}^2\right] \left(\frac{\delta\gamma}{\gamma_0} + \frac{\delta n}{n_0}\right) = c^2 \frac{\partial^2}{\partial z^2} \left(\frac{\delta\gamma}{\gamma_0} + \frac{\delta\gamma_F}{\gamma_{F0}}\right), \quad (17)$$

while relations between $\delta\gamma$ and $\delta\gamma_F$ obtained from Eqs. (6) and (7) are

$$\frac{\delta\gamma}{\gamma_0} = \alpha \left(\frac{\delta a}{a_0} - \frac{\delta\gamma_F}{\gamma_{F0}}\right) \\ \frac{\delta\gamma_F}{\gamma_{F0}} = \beta \frac{\delta n}{n_0} \quad (18)$$

with $\alpha = (1 - 1/\gamma_0^2)$ and $\beta = (1 - 1/\gamma_{F0}^2)/3$.

To obtain the dispersion equation, we assume that all perturbations depend on the coordinates and time like $\delta f \sim \exp(-i\omega t + ikz)$. After some algebra, Eqs. (15)–(18) lead to dispersion relation

$$\frac{\Omega_{e0}^2 A_a}{2} \frac{1}{D_e} (c^2 k^2 + \Omega_{e0}^2 - \omega^2) \left(\frac{1}{D_-} + \frac{1}{D_+}\right) = 1, \quad (19)$$

where $D_{\pm} = (\omega_0 \pm \omega)^2 - c^2(k_0 \pm k)^2 - \Omega_{e0}^2$ and $D_e = \omega^2 - \Omega_{e0}^2 - V_m^2 k^2 c^2$. Here, $V_m^2 = \beta(1 - \alpha)/(1 - \alpha\beta)$, $A_a = \alpha(1 - \beta)/(1 - \alpha\beta)$, which can be presented in the explicit form as

$$V_m^2 = \frac{1}{3} V_{F0}^2 \left(1 - \frac{a_0^2}{\gamma_m^2}\right) \left(1 - \frac{1}{3} V_{F0}^2 \frac{a_0^2}{\gamma_m^2}\right)^{-1}, \quad (20)$$

$$A_a = \frac{a_0^2}{\gamma_m^2 \gamma_{m0}^2}, \quad (21)$$

where V_{F0} is the dimensionless ($V_{F0} \rightarrow V_{F0}/c$) Fermi velocity defined by the relation $V_{F0} = (p_{F0}/\gamma_{F0})$ and $\gamma_{m0} = 1/\sqrt{1 - V_m^2}$.

The dispersion relation (19) has the form similar to the one obtained in Ref. 8. In this study, the authors investigated the SRS in classical relativistic plasmas where thermal effects were included by a waterbag model for electron distribution. It happens so that the description of a fully degenerate plasma is formally analogous to a waterbag model,¹⁶ and this analogy is reflected in the form of Eq. (19). Important differences, induced by plasma degeneracy, are contained in the expressions (20) and (21). Note; in a classical plasma, V_m is determined by the temperature (an independent parameter of the problem), while in a degenerate system, V_m depends on plasma density (determining the ‘‘Fermi temperature’’). For ultra-relativistic degeneracy ($R_0 \gg 1$), $V_m \rightarrow 1/\sqrt{3}$. Hence, the maximal value of the parameter $\gamma_{m0} \simeq 1.22$.

The solutions of Eq. (19) lead to resonance backward and forward SRS. Like in the classical plasma, under certain simplified assumptions, temporal growth rates ($\Gamma = \text{Im}(\omega)$) of instabilities can be obtained analytically. However, since the dispersion relation is a sixth order algebraic equation for ω , numerical solutions may be more useful.

The transverse EM and the plasma waves are coupled by the parameter $\chi = \Omega_{e0}^2 A_a$, and the growth rates of the above mentioned instabilities are proportional to a certain power of the coupling parameter χ . In the weakly degenerate case, $R_0 \ll 1$ and the finite strength of the field amplitude (a_0) $\gamma_m \approx \sqrt{1 + a_0^2}$, $V_m \ll 1$ and $\chi = \omega_{e0}^2/\gamma_m^3$. In this case, the dispersion relation (19) and, consequently, the results of instability coincide with the results obtained in the classical cold plasma embedded in the field of arbitrary strong radiation.^{6,7} Here, we would like to remark that our consideration is valid if the average kinetic energy of electrons ($\sim \epsilon_F$) is larger than their interaction energy ($\sim e^2 N_0^{1/3}$). This condition is fulfilled for a sufficiently dense plasma when $N_0 \gg 10^{23} \text{ cm}^{-3}$ ($1 \gg R_0 \gg 10^{-7}$).

The main features and peculiarities of SRS instability in the plasma with a finite level of degeneracy parameter R_0 can be deduced by analyzing Eqs. (20) and (21) for different limiting cases. We look for the solution of Eq. (19) in the form $\omega = \omega_L + \delta\omega$, where $\omega_L = \sqrt{\Omega_{e0}^2 + V_m^2 k^2 c^2}$. The maximum growth rates are obtained when the scattered wave is also resonant $D_-(\omega = \omega_L) = 0$, leading to the following relation:

$$\omega_0 = (\Omega_{e0}^2 + c^2 k^2 V_m^2)^{1/2} + ((k_0 - k)^2 + \Omega_{e0}^2)^{1/2}. \quad (22)$$

Using Eq. (22), one can show that in an underdense plasma ($\omega_0 > 2\Omega_{e0}$), the wave vector of resonant modes k lies in the range $0 < k < 2k_0$. The modes with $k > k_0$ and $k < k_0$ lead to backward and forward Raman instabilities, respectively. For a highly transparent plasma ($\omega_0 \gg \Omega_{e0}$), the backward Raman instability develops at $k \simeq 2\omega_0/c(1 + V_m)$ and $\omega_L = [\Omega_{e0}^2 + 4\omega_0^2 V_m^2/(1 + V_m)^2]^{1/2}$.⁸ Neglecting the

nonresonant term ($\sim 1/D_+$) in Eq. (19) and making approximations $D_- = -2(\omega_0 - \omega_L)\delta\omega$ and $\omega_L \gg \delta\omega$ (and recalling that $\Omega_{e0} = \omega_{e0}/\gamma_m^{1/2}$) for the growth rate, we get the following expression:

$$\Gamma = \frac{\omega_{e0} a_0}{\sqrt{2} \gamma_m^{3/2}} \frac{\omega_0(1 - V_m)}{\sqrt{(\omega_0 - \omega_L)\omega_L}}. \quad (23)$$

For forward Raman scattering instability $k \ll k_0 \approx \omega_0/c$, both downshifted and upshifted scattered waves are resonant modes $D_{\pm} = 2(\omega_L \pm \omega_0)\delta\omega$. The maximum growth rate is given as follows:

$$\Gamma = \frac{\omega_{e0}^2 a_0}{2\omega_0 \gamma_m^2 \gamma_{m0}} \quad (24)$$

occurs at $k = \Omega_{e0}\gamma_{m0}/c$ and $\omega_L = \Omega_{e0}\gamma_{m0}$.

As per expectations, in the limit of nonrelativistic degeneracy ($R_0 \ll 1$), Eqs. (23) and (24) tend to give the classical results for a cold plasma. For the finite level of degeneracy parameter R_0 and nonrelativistic strengths of the field ($a_0 \ll 1$), the growth rates of forward Raman instability is reduced by factor $(1 + R_0^{2/3})$ in comparison to the weakly degenerate case, while for backward instability, the reduction factor turns out to be $(1 + R_0^{2/3})^{3/4}$.

The most interesting regime, explored in this note, is that of a relativistic degenerate plasma ($R_0 \geq 1$) embedded in the field of a relativistically strong EM wave ($a_0 \gg 1$). In the regime of very strong radiation ($a_0/\gamma_0 \gg 1$ [see Fig. 1 and comments after Eq. (13)], even a highly degenerate plasma responds as a weakly degenerate one. For ultrarelativistic amplitudes, defined by $R_0/\gamma_0 < 1$, the Fermi velocity becomes small ($V_m \rightarrow 0$), and $\gamma_m \approx \sqrt{1 + a_0^2}$. Consequently, the character of the SRS instability, in particular its growth rate, tends to approach the cold classical plasma results.

In this letter, we investigated the linear stage of SRS instability of a propagating arbitrary amplitude circularly polarized EM wave in a degenerate electron plasma. For 1-D wave propagation, it is demonstrated that in the field of ultra-relativistic amplitude waves, the relativistic degenerate plasma effectively responds as a weakly degenerate plasma due to reduction of Fermi momentum. While for weak, non-relativistic amplitudes (but relativistic degenerate plasma), the growth rates of instability reduce with increasing density.

The elucidation of the SRS instability induced by ultra-strong EM waves in a relativistic degenerate plasma is highly pertinent to understanding the dynamics of X-ray pulses emanating from compact astrophysical objects. It will be equally relevant for an exploration of nonlinear interactions between intense laser pulses and a dense degenerate plasma; the latter class of physical systems are likely to be realized in the next-generation intense laser solid density plasma experiments.

This research was supported by the Shota Rustaveli National Science Foundation Grant (DI-2016-14). The research by GT was supported by the Knowledge Foundation at the Free University of Tbilisi.

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