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Electron acceleration by ponderomotive force in magnetized quantum plasma

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Abstract

The possibilities of electron acceleration by ponderomotive force of a circularly polarized laser pulse in magnetized quantum plasma have been explored. The basic mechanism involves acceleration of electron by the axial gradient in the ponderomotive potential of the laser. The quantum effects have been taken into account for a high-density plasma. The ponderomotive force of the laser is resonantly enhanced when Doppler up-shifted laser frequency equals the cyclotron frequency.

Keywords: Particle acceleration; Plasma-based accelerators; Ponderomotive force; Quantum plasma

1. INTRODUCTION

Electron acceleration by interaction of high-intensity laser pulse with plasma is of great interest because of its tremendous applicability in fundamental research and industrial use. There are several mechanisms for laser-driven electron acceleration, such as direct laser acceleration and acceleration by ponderomotive force of laser-induced in various forms (Tajima & Dawson, 1979; Khachatryan, 2002). The ponderomotive force accelerates the electron or excites a large amplitude plasma wave that can accelerate electrons indirectly. Many schemes of particle acceleration have been proposed since long-time back and recent developments are also there in the acceleration process. Optical mixing of laser light in a plasma and electron acceleration by relativistic electron plasma waves has been studied by Ebrahim in (1994). Pukhov (2004) investigated the particle acceleration process in relativistic laser channels; multi-GeV energy gain in a plasma wakefield accelerator and low-energy spread electron bunches (100 MeV–1 GeV) from laser wakefield acceleration has been studied by Hogan *et al.* (2005), Robinson (2006), etc. over the last decade. Kinetic modeling of intense short laser pulses propagating in plasmas as well as the ponderomotive effects on electron acceleration by plasma wave by short pulse lasers has been studied by Mora and Antonsen (1997) and Liu and Tripathi (2005). Ponderomotive acceleration of electron by a self-focused laser pulse has been studied by Singh *et al.* (2010). The interaction of arbitrarily polarized laser pulse with a tenuous plasma in the acceleration process has been studied by Sazergari *et al.* (2006). Numerical study has been carried out from much earlier times. Hartemann *et al.* (1995) studied the nonlinear ponderomotive scattering of relativistic electrons by intense laser field at focus in vacuum by employing one-dimensional (1D) plane wave pulse and 2D Gaussian pulses. Later, Pukhov and Vehn (1999) in their three-dimensional particle-in-cell (3D-PIC) simulation of intense short laser pulse interaction with plasma have observed the strong flow of relativistic electrons axially, moving with the laser pulse and generating 100 MG azimuthal magnetic field. The electron energies are far in excess of ponderomotive potential energy and the acceleration is a consequence of direct exchange of energy between electron and laser via betatron resonance phenomenon. Later in the same year Gahn (1999) experimentally observed generation of multi-MeV electron beam by direct laser acceleration in high-density plasma channels. Tsakiris *et al.* (2000) developed an analytical theory and a fully relativistic 3D single-particle code for direct laser acceleration of electron in radial electrical and azimuthal magnetic fields. Effect on an axial magnetic field and ion space charge on laser beat wave acceleration and a surfatron acceleration of electron has been studied by Prasad *et al.* (2009). Tanimoto (2003) studied the effect of self-induced azimuthal magnetic field on the direct electron acceleration by laser with stochastic phase disturbance.

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They found that apart from beam collimation electrons are accelerated to ultrahigh energies that are greater than the ponderomotive energy and that the acceleration is enhanced by increasing the strength of magnetic field. Liu *et al.* (2004) showed that the electron acceleration depends on laser intensity and the ratio of cyclotron frequency to laser frequency. Yu *et al.* (2002) examined similar configuration using linearly polarized laser. They have obtained electron acceleration to relativistic energies using weak magnetic field. The dependence of high-energy electron generation on the pulse duration of a high-intensity LFEX laser was experimentally investigated by Kojima (2016).

For plasma where the density is quite high and the de-Broglie thermal wavelength associated with the charge particle, that is, $\lambda_B = \hbar/2\pi m k_B T$ approaches the electron Fermi wavelength λ_{Fe} and exceeds the electron Debye radius λ_{De}(viz.,$\lambda_B \sim \lambda_{Fe} > \lambda_{De}$), the quantum effects become important. Furthermore, the quantum effects associated with the strong density correlation start playing a significant role when λ_B becomes of the same order or larger than the average inter-particle distance ($\sim n_0^{-1/3}$), that is, $n_0\lambda_B^3 \geq 1$ hold in degenerate plasma. However, the other condition for degeneracy is that the Fermi temperture T_F, which is related to the equilibrium density n_0 of the charged particles must be greater than the thermal temperture T of the system. The high-density, low-temperature quantum Fermi plasma is significantly different from the low-density, high-temperature 'classical plasma' obeying Maxwell–Boltzmann distribution. Over the last decade, there has been a growing interest in investigating new aspects of dense quantum plasmas by developing the quantum hydrodynamic (QHD) equations (Gardner & Ringhofer, 1996). The QHD equations are useful for studying numerous collective effects (Shukla & Eliasson, 2006) involving different quantum forces. Haas *et al.* (2000) presented a quantum multistream model by using a nonlinear system and derived the dispersion relations for one and two-stream plasma instability. Later, Anderson (2002) examined the statistical behavior of quantum plasma. Propagation of short-wavelength electromagnetic waves through magnetized quantum plasmas have been studied (Shokari et al., 2003; Shukla, 2006; Ali, 2006). The instabilities arising in quantum plasmas have also been studied (Ludin, 2007; Bret, 2008). Possibilities of magnetic field generation by ponderomotive force of electromagnetic waves in dense plasma has been explored (Shukla et al., 2010; Kumar & Tewari, 2012). Recently, it has been shown that the quantum effects can be important even in the classical regime (Opher et al., 2001; Brodin et al., 2008).

In the present work, we focus on the recently developed QHD model (Shukla & Eliasson, 2006; Haas *et al.* 2000). The QHD model consists of a set of equations describing the transport of charge density, momentum (including the Bohm potential) and energy in a charged particle system interacting through a self-consistent electrostatic potential. QHD is a macroscopic model and application is limited to those systems that are large compared with Fermi length of the species in the system. The advantages of the QHD model over kinetic descriptions are its numerical efficiency, the direct use of the macroscopic variables of interest such as momentum and energy and the easy way the boundary conditions are implemented.

In this paper, we have examined the possibilities of electron acceleration by a circularly polarized laser pulse in magnetized quantum plasma. The basic mechanism involves acceleration of electron by the axial gradient in the ponderomotive potential of the laser. The ponderomotive force of the laser is resonantly enhanced when Doppler-shifted laser frequency equals the cyclotron frequency. Such a study has not been reported in literatures so far.

2. THEORETICAL FORMULATION

Consider the propagation of a right circularly polarized laser pulse in quantum plasma in the direction of the static axial magnetic field $b_0\hat{z}$. The laser fields being (Sharma & Tripathi, 2009)

$$E_x = A\left(t - \frac{z}{\eta_g c}\right)\exp\left[-i\omega\left(t - \frac{\eta z}{c}\right)\right], \quad E_y = iE_x;$$

$$B_y = \left[\eta A + i\frac{(1-\eta\eta_g)}{\eta_g\omega}\frac{\partial A}{\partial t}\right]\exp\left[-i\omega\left(t - \frac{\eta z}{c}\right)\right], \quad B_x = iB_y;$$

where A for a Gaussian pulse is given by

$$A^2 = A_0^2\exp\left[\frac{-(t - z/\eta_g c - t_0)^2}{\tau^2}\right],$$

$\eta_g c = \eta c/[1 + \omega_p^2\omega_b/2\omega(\omega - \omega_b)^2]$ is the group velocity of laser pulse; τ is the laser pulse duration, $\eta = [1 - \omega_p^2/\omega(\omega - \omega_b)]^{1/2}$ is the refractive index of plasma; ω_p is the electron plasma frequency; $\omega_b = eb_0/mc$ is the cyclotron frequency; $-e$ is the electronic charge; m is the rest mass of electron; and c is the speed of light. The refractive index is crucially dependent on the static magnetic field, through the term $(\omega - \omega_b)$. This is due to the fact that the static magnetic field tends to rotate the electron. The laser field also rotates the electron, and hence at resonance ($\omega = \omega_b$), the two rotation frequencies resonate resulting in considerable enhancement in electron response and hence the refractive index.

The QHD equations (Jung, 2013; Li et al., 2014; Seadway, 2014; Wallin et al., 2014) governing the motion of electron in the presence of laser field and the static magnetic field are given by

$$\vec{F} = \frac{d\vec{p}}{dt} = -e\vec{E} - \frac{e}{\gamma mc}(\vec{p}\times\vec{B}) - \frac{v_F^2}{3n_0^2}\frac{\nabla n^3}{n} + \frac{\hbar^2}{2m_e}\nabla\left(\frac{1}{\sqrt{n}}\nabla^2\sqrt{n}\right) \tag{1}$$

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and the continuity equation

$$\frac{\partial n}{\partial t} + \nabla.\left(\frac{n\vec{p}}{\gamma m}\right) = 0, \quad (2)$$

where $n = (n_0 + n^{(1)})$ is the electron density, m is the electron's rest mass, $\hbar$ is the Planck's constant divided by 2π, and $v_F = [\hbar/m)(3\pi^2 n)^{1/3}]$ is the Fermi velocity. The third term on the right-hand side of Eq. (1) denotes the Fermi electron pressure ($P = mv_F^2 n^3/3n_0^3$). The fourth term is the quantum Bohm force and is due to the quantum corrections in the density fluctuation. The classical equation may be recovered in the limit of $\hbar = 0$. The ponderomotive force of the high-frequency laser pulse drives longitudinal waves with a frequency much smaller than ω, but fast enough for the dynamics to take place on the electron time scale. The ions form a neutralizing background in the dense plasma. Perturbatively expanding Eqs. (1) and (2) for first order of the electromagnetic field, we get

$$\frac{d\vec{p}^{(1)}}{dt} = -e\vec{E}^{(1)} - \frac{e}{\gamma_0 mc}\left[\vec{B}^{(0)} \times \vec{p}^{(1)}\right] - \frac{v_F^2}{n_0}\nabla n^{(1)} + \frac{\hbar^2}{4m^2}\left(\frac{1}{n_0}\nabla(\nabla^2 n^{(1)})\right), \quad (3a)$$

$$\frac{\partial n^{(1)}}{\partial t} + \frac{n_0}{\gamma_0 m}(\nabla.\vec{p}^{(1)}) = 0, \quad (3b)$$

where, $p^{(1)}$ is the quiver momentum. The fourth term of Eq. (3a) has been obtained by using perturbative expansion (Cao & Ren, 2008)

$$\nabla\left(\frac{1}{\sqrt{n}}\nabla^2\sqrt{n}\right) = \frac{1}{n_0}\left[\frac{1}{2}\nabla\nabla^2 n^{(1)} - \frac{1}{2n_0}\nabla n^{(1)}\nabla^2 n_0 - \frac{1}{4n_0}\nabla(2\nabla n_0.\nabla n^{(1)}) + \frac{1}{4n_0^2}n^{(1)}\nabla(\nabla n_0)^2 + \frac{1}{2n_0^2}(\nabla n_0)^2\nabla n^{(1)} + \frac{1}{n_0^2}(\nabla n_0.\nabla n^{(1)})\nabla n_0 - \frac{1}{n_0^3}(\nabla n_0)^2 n^{(1)}\nabla n_0\right].$$

Thus, the electron motion in the presence of laser and guiding magnetic fields is described by

$$\frac{dp_x^{(1)}}{dt} = -eE_x^{(1)} - \frac{eb_0}{\gamma_0 mc}p_y^{(1)} - \frac{e}{\gamma^{(1)}mc}p_z^{(1)}B_y^{(1)} - \frac{v_F^2}{n_0}\nabla n^{(1)} + \frac{\hbar^2}{4m^2}\left[\frac{1}{n_0}\nabla(\nabla^2 n^{(1)})\right], \quad (4a)$$

$$\frac{dp_y^{(1)}}{dt} = -eE_y^{(1)} + \frac{eb_0}{\gamma_0 mc}p_x^{(1)} - \frac{e}{\gamma^{(1)}mc}p_z^{(1)}B_x^{(1)} - \frac{v_F^2}{n_0}\nabla n^{(1)} + \frac{\hbar^2}{4m^2}\left[\frac{1}{n_0}\nabla(\nabla^2 n^{(1)})\right], \quad (4b)$$

$$\frac{dp_z^{(1)}}{dt} = -\frac{e(p_x^{(1)}B_y^{(1)} - p_y^{(1)}B_x^{(1)})}{\gamma^{(1)}mc} - \frac{v_F^2}{n_0}\nabla n^{(1)} + \frac{\hbar^2}{4m^2}\left[\frac{1}{n_0}\nabla(\nabla^2 n^{(1)})\right], \quad (4c)$$

where $\gamma^{(1)} = (1 + (p_x^{(1)2}/m^2c^2) + (p_y^{(1)2}/m^2c^2) + (p_z^{(1)2}/m^2c^2))^{1/2}$. Assuming the perturbed density to vary as $n^{(1)} = n_1\ (t - (z/\eta_g c))\exp\ [-i\omega\ (t - (\eta z/c))]$ and simultaneously solving Eqs. (4a) and (4b), we get the transverse momentum as

$$p_x^{(1)} = (\beta + n_1\beta_q)\left(t - \frac{z}{\eta_g c}\right)\exp\left[-i\omega\left(t - \frac{\eta z}{c}\right)\right], \quad (5a)$$

$$p_y^{(1)} = ip_x^{(1)}, \quad (5b)$$

where

$$\beta = -\frac{e(\partial A/\partial t)}{\omega_b^2[(\omega/\omega_b)(1 - (P_z\eta/\gamma^{(1)})) - (1/\gamma_0)]^2}\left[\left(1 - \frac{P_z\eta}{\gamma^{(1)}}\right)^2 - \frac{\omega_b P_z(1 - \eta\eta_g)}{\eta_g\omega\gamma^{(1)2}}\right] + \frac{eA(1 - (P_z\eta/\gamma^{(1)}))}{i\omega_b[(\omega/\omega_b)1 - (P_z\eta/\gamma^{(1)}) - (1/\gamma_0)]} + \frac{eA\omega_b(1 - (P_z/\eta_g\gamma^{(1)}))(1 - (P_z\eta/\gamma^{(1)}))}{\omega_b^3\gamma^{(1)2}[(\omega/\omega_b)1 - (P_z\eta/\gamma^{(1)}) - (1/\gamma_0)]^3}\frac{d\gamma^{(1)}}{dt},$$

$$\beta_q = \frac{eA(1 - (P_z\eta/\gamma^{(1)}))}{i\omega_b[(\omega/\omega_b)(1 - (P_z\eta/\gamma^{(1)})) - (1/\gamma_0)]}\left[\frac{v_F^2}{n_0 c}\left\{\frac{1}{\eta_g} - i\omega\eta\left(t - \frac{z}{\eta_g c}\right)\right\} + \frac{\hbar^2}{4m^2 n_0}\frac{\omega^2\eta^2}{c^3}\left\{\frac{3}{\eta_g} - i\left(t - \frac{z}{\eta_g c}\right)\right\}\right],$$

and

$$P_z = \frac{p_z^{(1)}}{mc}.$$

The perturbed electron density is obtained with the help of Eq. (3b) on substituting the relevant quantities

$$n^{(1)} = -\frac{(n_0(\beta + i\beta)/\gamma_0 m)(i\omega\eta\eta_g(t - (z/\eta_g c)) - 1)(t - (z/\eta_g c))}{\gamma_0\eta_g mc\left\{1 - i\omega(t - (z/\eta_g c)) + (n_0(\beta_q + i\beta_q)(i\omega\eta\eta_g(t - (z/\eta_g c)) - 1)/\gamma_0\eta_g mc)\right\}}\exp\left[-i\omega\left(t - \frac{\eta z}{c}\right)\right]. \quad (6)$$

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The energy exchange is governed by the relation,

$$\frac{d\gamma^{(2)}}{dt} = \frac{e}{m^2c^2\gamma_0}\{(\beta + n_1\beta_q) + i(\beta + n_1\beta_q)\}$$
$$A_0\left(t - \frac{z}{\eta_g c}\right)^2 \left\{\exp\left[-i\omega\left(t - \frac{\eta z}{c}\right)\right]\right\}^2. \quad (7)$$

From the above equation, it is clear that the energy exchange is critically dependent on the static magnetic field. The magnetic field tends to rotate the electrons about the line of force in the right-handed sense at the cyclotron frequency. The right circularly polarized wave also rotates the electrons in the clockwise direction at the frequency ω. Hence, resonance occurs at $\omega = \omega_b$ resulting in energy gain enhancement.

In our analysis, we consider the parameter values as in (Misra et al., 2009), that is, the particle density to be of the order of 10^{34} m^{-3}. Such a density regime is relevant for dense plasma environments (including outer layers of astrophysical plasmas, such as neutron stars and white dwarfs) (Ali et al., 2007). In this regime, the particle Fermi velocity becomes less than the speed of light c in vacuum and the Fermi screening length is greater than the average particle distance, a condition for the collective quantum effects in plasmas to be important. It has been reported that for intense laser–solid density plasma interaction experiments and for the next generation of laser-based plasma compression experiments, the electron number density n_0 can vary between 10^{30} and 10^{34} m^{-3} (Harding & Lai, 2006; Shukla & Eliasson, 2007; Sharma & Tripathi, 2009; Zhu & Ji, 2012) and v_F/c is in the range $0.005 \rightarrow 0.25$ (Wang et al., 2013).

Figure 1, shows the variation of normalized group velocity of the laser pulse $\eta_g$with normalized plasma frequency (ω_p/ω) for $\omega = 0.5\omega_b$. It is observed that the group velocity decreases with increase in plasma frequency. It is seen that η_g initially drops sharply with (ω_p/ω) , but afterwards this decrease is at a slower pace. In Fig. 2, the variation of normalized group velocity with (ω/ω_b) for $\omega_p/\omega = 0.3$ has been studied. It is observed that the group velocity $\eta_g \approx c$, except on and near resonance ($\omega = \omega_b$).

We have numerically solved Eq. (7) to investigate the electron dynamics for different values of plasma density, ratio of laser frequency to electron cyclotron frequency (ω/ω_b), normalized laser amplitude and the static magnetic field. In Figures 3 and 4, the electron energy γ as a function of $\xi(= t - \eta z/c)$ have been plotted for $A_0 = 0.2$ and 0.8, respectively. The energy gain is increased due to increase in normalized laser amplitude. Higher electron energy can be achieved by increasing the laser field amplitude. As far as the effect of plasma frequency on electron is concerned, both the refractive index and group velocity of the laser pulse are functions of plasma frequency. Their values (η and η_g) increase with decrease in ω_p. The energy gain of electron will be optimum near resonance. The variation of energy gain with normalized frequency (ω/ω_b) has been shown in Figure 5. The energy gain increases with increase in ω/ω_b. Near resonance, at $\omega/\omega_b \simeq 0.9$, the group velocity is found to be $\sim 0.20c$, while at $\omega/\omega_b \simeq 0.4$, the group velocity has a larger value ($\sim 0.94c$), even then the gain is larger near resonance. This is due to the fact that the magnetic field of the laser, which is responsible for the ponderomotive force is significantly reduced at lower values of group velocity $[\alpha(1 - \eta\eta_g)]$. Electron attains maximum acceleration near the normalized Doppler-shifted cyclotron resonance $([\omega(1 - P_z\eta/\gamma^{(1)})/\omega_b - 1/\gamma_0] \sim 0)$. Also at, $\omega/\omega_b \simeq 0.4$, $([\omega(1 - P_z\eta/\gamma^{(1)})/\omega_b - 1/\gamma_0] \sim 0.42)$ remains aways from the resonance values, that is, zero hence electron gain less energy. At resonance the energy gain becomes infinity and the theory breaks down. After resonance, for $\omega/\omega_b > 1$, the term $[\omega(1 - P_z\eta/\gamma^{(1)})/\omega_b - 1/\gamma_0]$ increases and consequently the electron can gain higher energy.

Figure 6 depicts the variation of electron energy with ξ for $A_0 = 0.6$. The dashed line shows the variation as per the QHD model, while the solid line denotes the trend for a

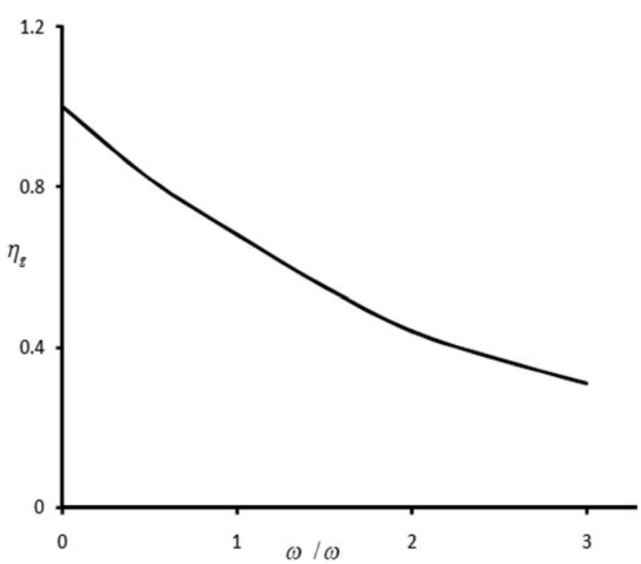

Fig. 1. Variation of η_g with ω_p/ω for $\omega = 0.5\omega_b$.

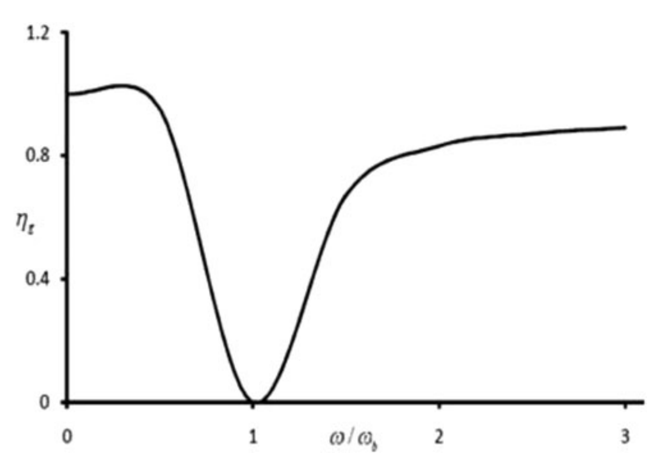

Fig. 2. Variation of η_g with ω/ω_b for $\omega_p = 0.3\omega$.

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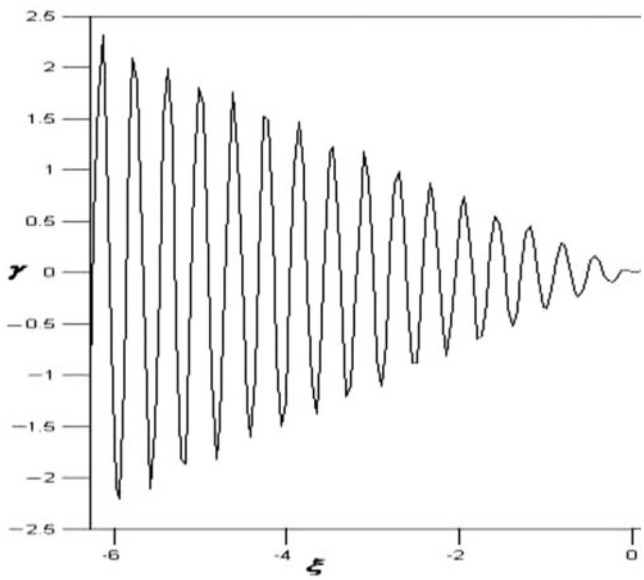

Fig. 3. Variation of γ with ξ for $A_0 = 0.2$, $n_0 = 10^{34}$ m^{-3}, $\omega_p = 0.3\omega$, $\omega = 0.5\omega_b$.

non-QHD setting (at $\hbar = 0$). It is found that the quantum effects damp the electron energy by about 14%. This damping is due to the diffraction induced by quantum effects.

3. DISCUSSION

A detailed theory for particle acceleration by the non-stationary ponderomotive force of a large-amplitude circularly polarized electromagnetic wave in a very dense, magnetized quantum plasma has been presented. The effects

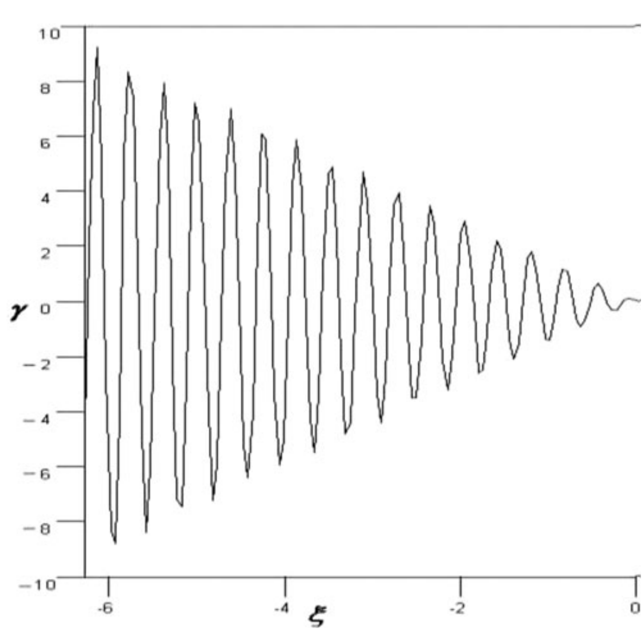

Fig. 4. Variation of γ with ξ for $A_0 = 0.8$, $n_0 = 10^{34}$ m^{-3}, $\omega_p = 0.3\omega$, $\omega = 0.5\omega_b$.

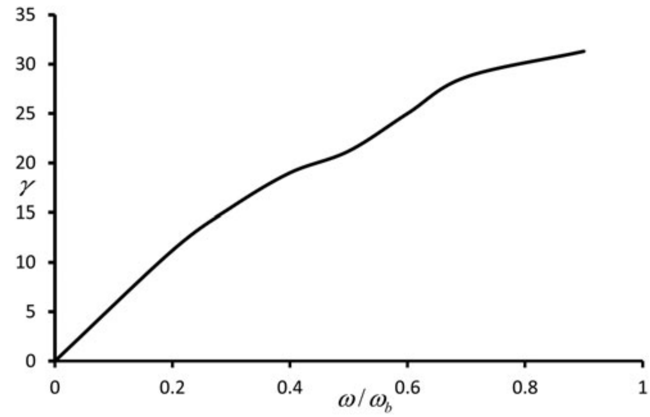

Fig. 5. Variation of γ with ω/ω_b for $A_0 = 0.2$, $z = -6$, $n_0 = 10^{34}$ m^{-3}.

associated with the Fermi pressure and the Bohm potential have been incorporated. The ponderomotive acceleration of electrons by a Gaussian laser pulse is significantly affected by the presence of a magnetic field. An electron gains energy during the rising part of the laser pulse and loses during the trailing part. The quantum diffraction effects also play a crucial role by modifying the energy exchange rate. The non-stationary radiation pressure creates a slowely varying electric fields and current, which contributes to ponderomotive acceleration. The non-oscillatory quantum terms are embedded in β_q. When $\omega < \omega_b$, the first-order velocity due to the laser is very small as the resonance condition is not satisfied and hence the ponderomotive force produced is also insignificant. For $\omega > \omega_b$, the pondermotive force produced is significant and considerable acceleration of electrons can be obtained in this regime. In practical applications, the ponderomotively accelerated ultrahigh energy electrons in the rising part of the pulse can be easily extracted by impinging the pulse into a solid target or through an overdense plasma seperator (Yu et al., 2000; Miyauchi et al., 2004). The pulse will then be reflected and the high-energy electrons are released into the target without suffering any deceleration. The electron energy is reduced by nearly 8% due to the

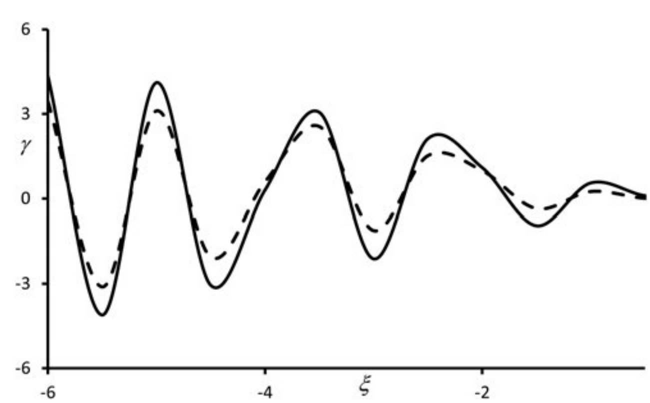

Fig. 6. Variation of γ with ξ for $A_0 = 0.6$, $n_0 = 10^{34}$ m^{-3}, $\omega_p = 0.3\omega$, $\omega = 0.5\omega_b$.

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quantum diffraction effects. This can be compensated by further increasing the strength of the applied magnetic field. The present study has relevance to the environment of dense astrophysical plasmas, quantum free-electron lasers, as well as to the next generation of intense laser–solid density plasma interaction experiments.

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