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Linear theory of quantum two-stream instability in a magnetized plasma with a transverse wiggler magnetic field

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Abstract

A fluid description is used to study the properties of two-stream instability due to interaction of a non-relativistic electron beam with quantum magnetized plasma and transverse wiggler magnetic field. It is assumed that the background plasma provides charge and current neutralization of the electron beam. The dispersion relation is obtained by solving and linearizing fluid-Maxwell equations. The resulting dispersion equation is analyzed numerically over a wide range of system parameters. The results of quantum and classical treatments are compared numerically, with including the effects of wiggler on the dispersion relation. It is found that the transverse wiggler magnetic field can strongly improve the instability of quantum plasma as well as classical plasma.

Keywords: Quantum plasma; Transverse wiggler magnetic field; Two-stream instability

1. INTRODUCTION

The interaction of electron beam with plasma plays a key role in various fields of physics and the theoretical study of the linear regime of beam-plasma instabilities has loomed large within most plasma physics textbooks. When an electron beam passes through a plasma, it quickly induces a return current. This current which fully neutralizes the beam current is carried by the plasma electrons. The resulting system undergoes various instabilities, multi-stream, two-stream, filamenation, and Weibel, that can disrupt beam propagation (Weibel, 1959; Fried, 1959; Davidson *et al*., 1975; Bret, 2009; Michno & Schlickeiser, 2010). The two-stream parallel unstable mode is found when the wave vector becomes aligned with the direction of beam propagation, while for the filamentation and Weibel modes the wave vector aligned perpendicular with the streams. In rather high density relativistic beams, filamentation modes grow fastest. On the other hand, for Maxwellian electron beam-plasma system in the non-relativistic regime, two-stream modes grow faster than any other unstable mode. These two instabilities are beam based which means they need a beam to exist. Two-stream instabilities are traditionally known to be important for the electron heating in intense laser-plasma interaction experiments (Thode & Sudan, 1973) as well as in astrophysical relativistic shocks (Nakar *et al*., 2011), and could be important for pulsar glitches (Andersson *et al*., 2003; Samuelsson *et al*., 2010), where super fluid neutrons and super conducting protons co-exist with relativistic electrons (Andersson *et al*., 2004).

On the other hand, the topic of quantum plasmas has received considerable attention in the recent advances. The main reason for this is that the quantum plasma has wide ranging applications in solid state physics, nano-scale objects such as nano-wires, quantum dots (Shpatakovskaya, 2006; Wei *et al*., 2007; Ang *et al*., 2006), ultrasmall electronic devices in microelectronics (Markowich *et al*., 1990), dense astrophysical plasmas (Chabrier *et al*., 2002; Opher *et al*., 2001; Jung, 2001), laser fusion plasmas(Glenzer *et al*., 2007), and in next-generation high intensity light sources (Marklund & Shukla, 2007; Mourou *et al*., 2006). It is well known that the quantum plasmas are characterized by their low temperatures and high densities in sharp contrast to the high temperatures and low densities of the classical plasmas (Shukla, 2006; Shukla *et al*., 2006; Haas *et al*., 2003*a*). The quantum fluid equations have been further used to describe the properties of the instabilities associated with the three-stream quantum plasmas (Haas *et al*., 2003*b*), the quantum dusty plasmas (Ali & Shukla, 2007), the electron-positron-ion quantum plasmas (Mushtaq & Khan, 2008), and the magnetized quantum multi-stream (Ren *et al*., 2008). The hydrodynamic formalism has also been used to investigate the quantum filamentation instability in field configurations with and

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without axial guide magnetic field (Bret, 2008; 2007; Hasanbeigi *et al.*, 2012; 2014; Mehdian *et al.*, 2013*a*; 2013*b*).

Although, many researcher have aseassed the electron beam plasma instability, to the author's knowledge the effects of the transverse helical magnetic field on the quantum two-stream instability have not been considered yet. The main purpose of this paper is to show how the quantum two-stream instability can be affected by external helical wiggler magnetic field. Due to importance of wiggler magnetic field, especially in free electron lasers and accelerators, we devote the present research to the evaluation of quantum effects on the two-stream instability in the presence of helical wiggler. The wiggler is used to induce a perpendicular wave motion in the electron beams. Here the direction of the wiggler magnetic field is perpendicular to the beam propagation direction. A cold fluid model for a cold electron beam passing through a cold quantum magnetized plasma, accounting for a wiggler magnetic field and a return electronic current. The plasma is assumed to be cold so that before the passage of the electron beam, the plasma electrons are at rest and the magnetic field does not affect them. Collisions as well as thermal motion of the beam and plasma electrons are also neglected. The layout of this paper is organized as follows. A description of the physical model and the equilibrium configuration are presented in Section 2. The system of equations that govern the motion of a non-relativistic electron beam through a background plasma and external magnetic field are also derived in this section. The full three-dimensional dispersion relation of two-stream instability is obtained in Section 3. Finally, Section 4 is devoted to the numerical results and discussion.

2. EQUILIBRIUM CONFIGURATION OF SYSTEM AND FIRST EQUATIONS

The equilibrium configuration illustrated in Figure 1 corresponds to a cold non-relativistic electron beam of density n_{b0} and velocity v_b propagating through background magnetized plasma in a configuration which consists of a helical wiggler, $\mathbf{B}_w = B_w(\hat{e}_x \cos k_w z + \hat{e}_y \sin k_w z)$, and a uniform axial solenoidal field, $B_0 e_z$. Here B_w and k_w are the magnitude of the helical magnetic field and wave number, respectively. When the beam enters the magnetized plasma, it is deflected by $\mathbf{B}_w$. The solutions of the electron orbits for this configuration can be obtained as

$$\mathbf{v}_{0b} = \frac{v_b \Omega_w}{\Omega_0 - k_w v_b}(\cos(k_w z)\mathbf{e}_x + \sin(k_w z)\mathbf{e}_y) + v_b \mathbf{e}_z, \quad (1)$$

where $\Omega_{0,w} \equiv q B_{0,w}/\text{mc}$, and the symbol q and m are used to denote the magnitude of electron charge and electron rest mass, respectively. The beam density and current are assumed to be sufficiently low, so that the effects of equilibrium self-electric and self-magnetic fields can be neglected (Davidson, 2001). On the other hand, a sufficiently dense plasma can also neutralize the injected current. Due to the changing magnetic flux of a propagating bunch, return current is induced in the plasma. Therefore, the current carried by the injected beam electrons is neutralized by a return current in the plasma. The plasma electrons have the density and drift velocity n_{p0} and $\mathbf{v}_{p0}$, respectively, so that $n_{b0}\ \mathbf{v}_{b0} = -n_{p0}\ \mathbf{v}_{p0}$. The ions form a motionless neutralizing background represented by a positive charge density $n_i = (n_{p0} + n_{b0})/Z$, where n_{p0}, v_{p0}, and Z are plasma electron density, plasma electron velocity and ion charge, respectively. The above imposed relations imply that the plasma system is both charge and current neutralized.

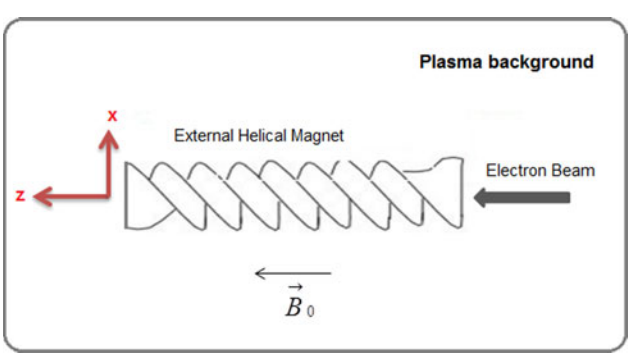

Fig. 1. (Color online) Schematic diagram of a beam plasma system to be adopted in the present paper.

3. DISPERSION EQUATION

To study the propagation of wave in this system, the Maxwell equations will be used along with the the cold electron momentum transfer equation,

$$\frac{\partial \mathbf{v}_j}{\partial t} = -\frac{q}{m}(\mathbf{E} + \frac{\mathbf{v}_j \times \mathbf{B}}{c}) + \frac{\hbar^2}{2m^2}\nabla(\frac{\nabla^2\sqrt{n_j}}{\sqrt{n_j}}), \quad (2)$$

and equation of continuity,

$$\frac{\partial n_j}{\partial t} + \nabla.(n_j \mathbf{v}_j) = 0, \quad (3)$$

where the index j for the beam and plasma component is b and p, respectively. Quantum corrections are clearly contained within the so-called Bohm potential by means of the $\hbar^2$ term in the above equation, which can be obtained from the momentums of the non-relativistic Wigner function. In order to develop a first order perturbation theory, the electron number density n_j, the electron velocity v_j, the electric field $\mathbf{E}$, and the magnetic field $\mathbf{B}$ will be expressed in the form

$$\begin{aligned} n &= n_{j0} + n_{j1}, \\ \mathbf{v}_j &= \mathbf{v}_{j0} + \mathbf{v}_{j1}, \\ \mathbf{E} &= \mathbf{E}_1, \\ \mathbf{B} &= \mathbf{B}_0 + \mathbf{B}_1, \end{aligned} \quad (4)$$

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where the unperturbed magnetic field $\mathbf{B}_0 = B_0\,\hat{e}_z + \mathbf{B}_w$ consists of the transverse helical wiggler and the axial field of magnitude B_0. We express the axial and time dependence of all perturbed parameters by application of Floquet's theorem (Sadegzadeh *et al.*, 2012; Hwang *et al.*, 2002; Jha & Kumar, 1998) for periodic systems in the general form $F = \sum_{n=-\infty}^{+\infty} f_n e^{i(k_n z-\omega t)}$, where k_n and ω are the wave numbers and frequency of wave and $k_n = k + nk_w$. As a consequence, the linearized continuity and momentum transverse equations yeild

$$v_{jxn} = -\frac{ei}{m\Omega_{jn}}\Big(-\frac{B_w(v_{zn+1}-v_{zn-1})}{2ci}+\frac{B_0 v_{jyn}}{c} + E_{xn}\Big(1-\frac{\beta_0 ck_n}{\omega}\Big)\Big), \tag{5}$$

$$v_{jyn} = -\frac{ei}{m\Omega_{jn}}\Big(-\frac{B_0 v_{jxn}}{c}+\frac{B_w(v_{zn-1}+v_{zn+1})}{2c} + E_{yn}\Big(1-\frac{\beta_0 ck_n}{\omega}\Big)\Big), \tag{6}$$

$$v_{jzn} = -\frac{ei}{m\Omega_{jn}}\Big[\frac{B_w(v_{xn+1}-v_{xn-1})}{2ci}-\frac{B_w(v_{yn-1}+v_{yn+1})}{2c} + \frac{S_j(k_{n+1}E_{yn+1}-k_{n-1}E_{yn-1})}{2i\omega} + \frac{S_j(k_{n-1}E_{xn-1}+k_{n+1}E_{xn+1})}{2\omega}\Big]+\frac{\hbar^2 n_{jn}k_n^3}{4m^2 n_{0j}\Omega_{jn}}, \tag{7}$$

$$n_{1j} = \frac{n_{0j}k_n}{\omega - v_{jz0}k_n}v_{jzn} \tag{8}$$

where

$$S_j = \frac{v_{jz0}\Omega_w}{\Omega_0 - k_w v_{jz0}},\ \Omega_{jn} = \omega - v_{jz0}k_n, \tag{9}$$

and $\mathbf{B}_1$ has been eliminated in momentum transverse equations by use of $\mathbf{B}_1 = c/\omega \mathbf{k}\times\mathbf{E}_1$. Making use of Eqs. (5)–(8), the linearized current density, $\Sigma_j(n_{1j}\,\mathbf{v}_{j0} + n_{j0}\,\mathbf{v}_{j1})$, component amplitudes are given by

$$J_{xn} = -e\Big[n_{j0}v_{jxn}+\frac{1}{2}n_0 S_j\Big(\frac{k_{n-1}v_{jzn-1}}{\Omega_{jn-1}}+\frac{k_{n+1}v_{jzn+1}}{\Omega_{jn+1}}\Big)\Big],$$
$$J_{yn} = -e\Big[\frac{n_0 S_j\Big(\frac{k_{n+1}v_{jzn+1}}{\Omega_{jn+1}}-\frac{k_{n-1}v_{jzn-1}}{\Omega_{jn-1}}\Big)}{2i}+n_{j0}v_{jny}\Big], \tag{10}$$
$$J_{zn} = -ev_{jzn}n_{j0}\Big(\frac{v_{0jz}k_n}{\Omega_{jn}}+1\Big).$$

Substituting Eqs. (5)–(10) into the wave equation, $\nabla\times(\nabla\times\mathbf{E}_1) + c^{-2}\partial^2\ \mathbf{E}_1/\partial t^2 + 4\pi c^{-2}\partial\mathbf{J}_1/\partial t = 0$, then gives a dispersion equation with the form $det\,\overleftrightarrow{\mathbf{D}} = 0$ with

$$\overleftrightarrow{\mathbf{D}} = \begin{pmatrix} D_{11} & D_{12} & 0 \\ D_{12}* & D_{22} & 0 \\ 0 & 0 & D_{33} \end{pmatrix}, \tag{11}$$

where the superscript * refers to the complex conjugate and the elements of the dispersion tensor $\overleftrightarrow{\mathbf{D}}$ are complicated functions of the system parameters. The dispersion equation is a polynomial in x and can be simplified to the form $\sum_{n=0}^{20} a_n x^n = 0$, where the coefficients a_n are very complicated, $x = \omega/\omega_p$ and $\omega_p = (\frac{4\pi n_{p0}q^2}{m})^{1/2}$. It should be noted that in the limit of zero magnetic field, $\Omega_{0,w}\to 0$, the dispersion equation is the same as the one found by Hass *et al.* (2000; 2009). The dispersion properties of cold quantum magnetized plasma in the presence of the transverse wiggler magnetic field can be described by analyzing the 20° polynomial in ω. Owning to complicated nature of the dispersion relation, we will only treat it numerically.

4. NUMERICAL RESULTS AND DISCUSSIONS

A numerical study has been made to illustrate two-stream instability of quantum magnetized plasma with transverse wiggler magnetic field. Results of numerical solutions of the 20° polynomial dispersion equation for the complex normalized wave frequency, x, are described below. Classical and quantum limit correspond to the case $\Theta = (\frac{\hbar\omega_p}{2mc^2})^2 = 0$ and $\Theta = 9\times10^{-6}$, respectively. First the results of quantum and classical treatments will be compared numerically, including the

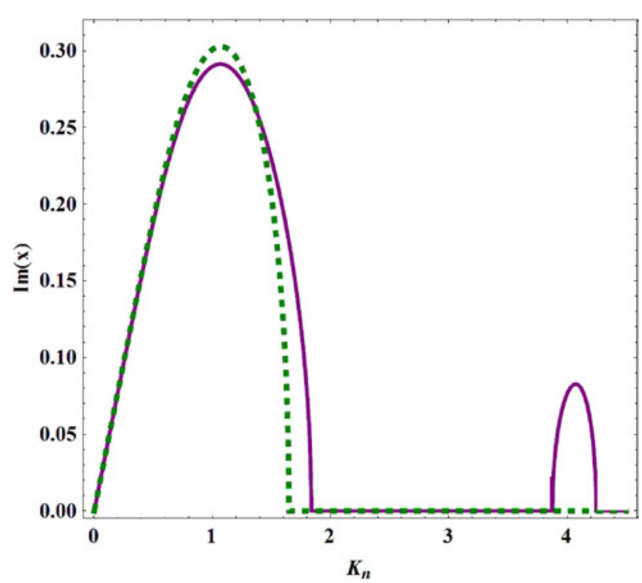

Fig. 2. (Color online) Plot of normalized growth rate in term of normalized wavenumber K_n for (**a**) quantum magnetized plasma, (**b**) classical magnetized plasma, and (**c**) classical non-magnetized plasma. The parameters are $\alpha = 0.1$, $\beta = 0.1$, $\Omega = \Omega_w = 3.51\times10^{11}\,Hz$, $\Theta = 9\times10^{-6}$.

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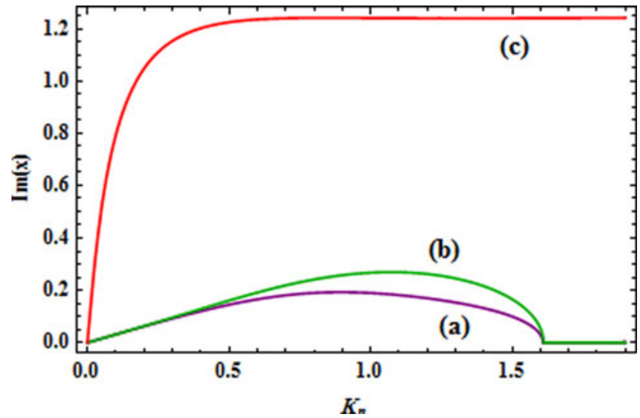

Fig. 3. (Color online) Graph of normalized growth rate versus normalized wavenumber K_n for quantum limit (the solid curve) and classical limit (the dash curve). The parameters are: $\alpha = 0.15$, $\beta = 0.15$,. The other parameters are the same as in Figure 2.

effects of wiggler on the dispersion relation. The normalized growth rate $Im(x)$ as a function of the normalized wave number $K = \frac{k_n v_b}{\omega_p}$ is shown in Figure 2 for (a) quantum magnetized plasma, (b) classical magnetized plasma, and (c) classical plasma. Here Ω_w ,$\alpha = n_{b0}\ /n_{p0}$,k_w ,Ω, and $\beta = v_b\ /c$ were taken to be 3.51×10^{11} Hz, 0.1, $2\pi/3\ cm^{-1}$, 3.51×10^{11} Hz, and 0.1, respectively. It is easy to see that the increasing the wave number of classical plasma, $\Theta = \Omega = 0$, results in a higher wave growth rate until it reaches a saturation value. The behavior of magnetized plasma is completely different compared to the non-magnetized plasma case. This is evident in the figure since, for example, the growth rate increases with increasing K, achieves a maximum value at $K = K_{max}$ and then decreases to zero with K for its values within a higher range, $K > K_{max}$. It is also seen from this figure that there is no saturation for classical magnetized plasma and its growth rate is grater than that of for quantum magnetized plasma. The basic differences between the growth rate of a quantum and a classical magnetized plasma are illustrated in Figure 3. The dashed and solid curves in this figure correspond to the classical and quantum magnetized plasma, respectively. For quantum magnetized plasma, two peaks are observed. The larger peak is occurring in long wave lengths whereas the smaller peak corresponding to short wave lengths. The smaller peak is not seen in the classical regimes (Bret & Hass, 2010). Figure 4 shows the variation of the growth rate with K in a quantum plasma in

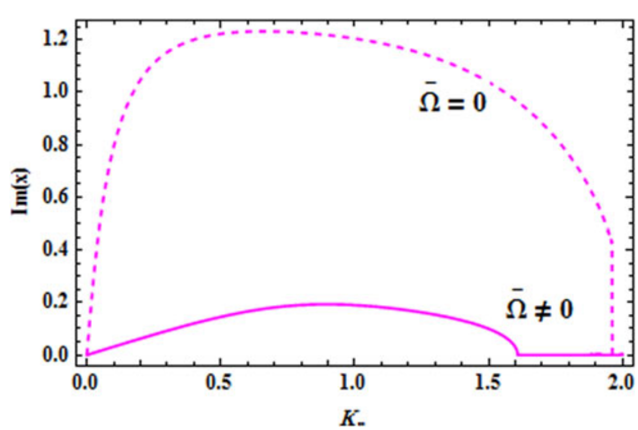

Fig. 4. (Color online) Graph of normalized growth rate as a function of normalized wavenumber for quantum plasma with and without axial magnetic field. Parameters are the same as in Figure 2.

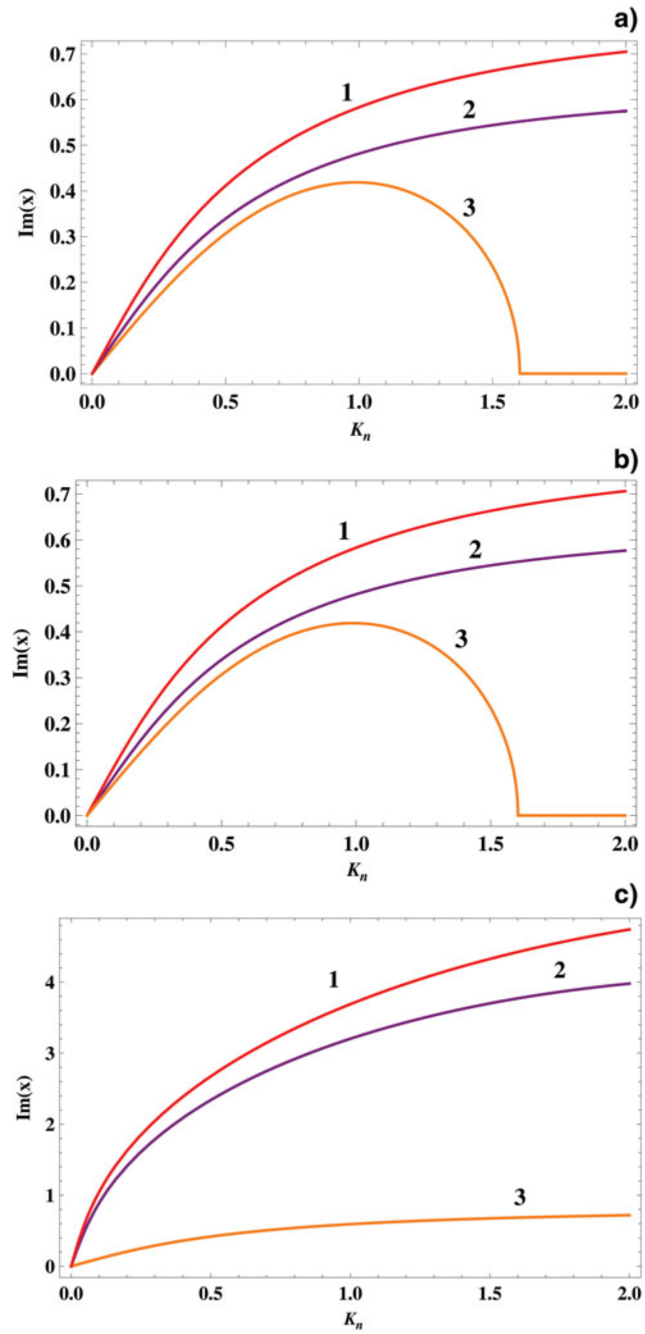

Fig. 5. (Color online) Effect of external periodic magnetic field on the growth rate of magnetized plasma for (**a**) quantum limit, (**b**) classical limit and (**c**) classical non-magnetized plasma. The curves 1, 2, and 3 correspond to the case where $\Omega_w = 8.7 \times 10^{1}$ 1 Hz, 7.02×10^{1} 1 Hz, and 0.8×10^{1} 1 Hz. The parameters are $\alpha = 0.5$, $\beta = 0.5$. The other parameters are the same as in Figure 2.

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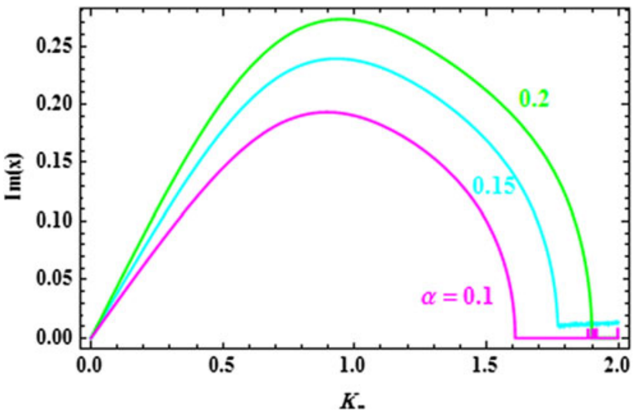

Fig. 6. (Color online) Plots of growth rate of quantum magnetized plasma versus K_n for for $\beta = 0.1$, $\Omega = \Omega_w = 3.51 \times 10^{11}\,Hz$, $\Theta = 9 \times 10^{-6}$ and several values of α.

the presence of the wiggler magnetic field. When Ω_w, α, Θ, and β are assumed to be constant, Figure 4 shows the effect of axial magnetic field on the growth rate of quantum plasma. It is obvious from this figure that the growth rate and the instability bandwidth are increased due to a decrease in the axial magnetic field.

To further understand the effects of wiggler on properties of the growth rate, we compare in Figure 5 the growth rate of instability in the presence of wiggler for different values of wiggler frequency. Shown in Figure 5 are plots of the growth rate versus K for three cases corresponding to $\Omega_w = 8.7 \times 10^{11}\,Hz$ (curve 1), $7.02 \times 10^{11}\,Hz$ (curve 2), and $\Omega_w = 0.8 \times 10^{11}$ (curve 3). It is evident from the Figure 5 that the growth rate is improved substantially by the presence of the wiggler magnetic field and the instability bandwidth becomes wider. The growth rate of instability can be affected significantly by the ratio of electron beam density to the plasma density, α. This is illustrated in Figure 6 where the normalized maximum growth rate of quantum plasma, $Im(x)$, is plotted versus the density ration α. Each curve corresponds to a fixed value of α. Note from the figure that the increased density ration, α, will increase the growth rate of instability and cut-off wave number. Therefore, the bandwidth of the instability is quite broad at large values of α.

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