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Influence of varying magnetic field on nonlinear wave excitations in collisional quantum plasmas

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Abstract: The nonlinear wave excitations arising from the spatially varying magnetic field in the quantum plasma environment are investigated in the frame work of quantum hydrodynamic model. In the weakly nonlinear, dispersive and dissipative limit it is shown that the varying magnetic field and collision-induced excitations can be described by a modified form of Korteweg-de Vries–Burgers' type model equation. It is found that the dissipation terms (Burgers' and collisional term) arise due to spatially varying magnetic field and the ion-neutral collisions. The numerical solutions of this equation predict that the localized soliton solutions decay algebraically due to the combined effect of varying magnetic field and collision by radiating oscillatory pulses behind the propagating soliton.

Keywords: dissipative quantum plasmas; electrostatic waves; nonlinear phenomena.

1 Introduction

The nonlinear propagation of low frequency electrostatic waves in magnetized quantum plasmas has recently gained much interest in dense astrophysical environments [1]. In dense quantum plasmas, degenerate electrons pursue Fermi pressure law, and there are typically quantum forces associated with the Bohm potential, which generate wave dispersion at nanoscales [2]. In the last several years, several researchers made their efforts on dense quantum plasmas due to its novelty as well as potential application prospects in ultrasmall electronic devices [3], in dense astrophysical objects like white dwarfs, neutron stars [4–9] and also in intense laser-solid interaction experiments [10]. In metallic nanostructures the quantum diffraction effect of charged particles has a significant role in the collective processes in plasmas. Spectral measurements of X-ray Thomson scattering [11] experimentally verify the influence of the quantum diffraction on the dispersive character of electrostatic waves in degenerate plasmas.

Most of the astrophysical systems comprise of magnetic field. Moreover, in space or laboratory plasma, motion of plasma particles generates a magnetic field. So, the fundamental characteristics of plasma motion in presence of magnetic field depend vitally on our ability to describe nonlinear behavior of magnetized plasma. Linear and nonlinear acoustic waves have been investigated in quantum plasmas by considering the quantum mechanical effects of plasma particles in different plasma media. For example, the behavior of low temperature and high density quantum plasmas was first studied by Pines [12]. Haas *et al.* [13] derived the dispersion relation for the two-stream instability and identified a new pure quantum branch to describe the dynamics of quantum plasmas with the help of nonlinear Schrödinger–Poisson model. Based on the Wigner–Poisson dynamical equations, Anderson *et al.* [14] described a statistical multistream model of quantum plasmas. Haas *et al.* [15] used the quantum hydrodynamic (QHD) model to study the linear and nonlinear quantum ion-acoustic waves (IAWs) in the limit of small mass ratio of charge carriers. Within the framework of the QHD approximation, the quantum effects on linear and nonlinear propagation of electrostatic waves have been investigated in the past by various authors [16–22]. Recently, Roy *et al.* [23] investigated face-to-face interaction between multisolitons in a fermionic quantum plasma. Singh *et al.* [24] addressed the problem of heavy nucleus-acoustic excitations in magnetorotating quantum plasma by deriving Zakharov–Kuznetsov–Burgers equation. More recently, Saha *et al.* [25] studied the dynamical properties of electrostatic IAWs in a dense Thomas–Fermi magnetoplasma.

The existence of strong magnetic fields in white dwarfs was predicted by Ginzburg [26]. The magnetic field plays a significant role in the study of the neutron star atmospheres and their radiation [27]. Haas [28] formulated a quantum magnetohydrodynamic (QMHD) model to establish the equilibrium conditions and the importance of

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magnetic field for dense astrophysical plasmas, such as, the interior of massive white dwarfs or the atmosphere of neutron stars. The linear and nonlinear behavior of the slow and fast magnetosonic modes were obtained using QMHD model that includes the Bohm and spin terms [29]. By using the QMHD model, Kaur *et al.* [30] investigated the nonlinear ion-acoustic (IA) cnoidal waves in magnetized quantum plasmas. Asenjo [31] discussed the effect of quantum corrections on the propagation of low frequency magnetosonic waves with mobile electrons and ions by considering the Bohm potential and the spin magnetization energy of electrons. In dense plasmas, the equilibrium density of charged carriers and the static ambient magnetic field can be nonuniform with finite scale lengths. Shukla and Stenflo [32] described the existence of new drift modes in a nonuniform quantum magnetoplasma and observed that the electron quantum correction significantly modified the electron drift-wave frequency. Moslem *et al.* [33] investigated electrostatic waves in a nonuniform quantum magnetoplasma and reported that the effects of quantum parameters, velocity, density and magnetic field inhomogeneities give rise to both oscillatory and purely growing instability in the local approximation regime. Misra [34] studied the amplitude modulation of low-frequency electrostatic drift-wave envelopes in a nonuniform quantum magnetoplasma.

In all the earlier investigations, external magnetic field has been taken constant. However, spatially varying magnetic field take place in the plasma transport processes both in laboratory as well as in different space plasma environment [35, 36]. A varying magnetic field causes changes in the dynamics of the electrons and ions in a dense plasma environment. It affects the ionization balance and the plasma spatial distribution. Transport properties of two-dimensional electron gas (2DEG) under a periodic magnetic field were theoretically investigated by Xue and Xiao [37]. Peeters and Vasilopoulos [38] studied quantum transport of a 2DEG in presence of a spatially modulated magnetic field. Wu and Ulloa [39] investigated electronic states and collective excitations of a 2DEG in presence of a spatially modulated magnetic field. Moreover, Loukopoulos and Tzirtzilakis [40] studied the problem of the biomagnetic fluid flow in a channel under the influence of an applied spatially varying magnetic field. Recently, Pakzad *et al.* [41] studied the behavior of small amplitude IA solitary waves in a nonrelativistic framework for classical plasmas under the influence of a spatially varying magnetic field. They showed that, in the presence of a varying magnetic field, solitons radiate some amount of energy during their propagation through the varying magnetic field, and the radiated energy emerges as backward moving shock waves. However, to the best of our knowledge, no attempt has been made to study the nonlinear propagation of quantum IAWs in a collisional quantum plasma in presence of a spatially varying magnetic field. Motivated by such considerations, we derive the modified form of Korteweg-de Vries–Burgers (KdVB) type equation to describe the nonlinear quantum IAWs excited by spatially varying magnetic field in a collisional plasma and discuss the potential utility of propagating waves influenced by several plasma parameters. The organization of the paper is as follows. In Section 2, we derive the nonlinear evolution equation that governed the dynamics of the nonlinear waves within the framework of QMHD model. In Section 3, we discuss the numerical solutions of the nonlinear equation with its graphical representations. Finally, a brief summary of the results is provided in Section 4.

2 Theoretical model and derivation of evolution equation

We consider the nonlinear propagation of electrostatic IAWs in a collisional quantum magnetoplasma consisting of positively charged inertial nondegenerate ions and inertialess degenerate electrons. At equilibrium, both electrons and ions have equal number density, say n_0. We introduce a space–dependent slowly varying external magnetic field $\mathbf{B} = B(\mathbf{r})\hat{z}$, i.e., the magnetic field in our model is a function of space [41]. The ion-neutral collisions are taken into account in the momentum conservation equations for ions as a simple relaxation term [42]. Thus, in collisional magnetized quantum plasma the basic set of QHD equations are [43, 44]

$$\frac{\partial n_i}{\partial t} + \nabla.(n_i\mathbf{v}) = 0, \tag{1}$$

$$\frac{dv_i}{dt} = -\frac{e}{m_i}\nabla\phi + \frac{e}{m_i}(\mathbf{v_i}\times\mathbf{B}) - \gamma_i\mathbf{v_i}, \tag{2}$$

$$0 = \frac{e}{m_e}\nabla\phi - \frac{\nabla p_e}{m_e n_e} + \frac{\hbar^2}{2m_e}\nabla\left(\frac{\nabla^2\sqrt{n_e}}{\sqrt{n_e}}\right), \tag{3}$$

$$\nabla^2\phi = \frac{e}{\epsilon_0}(n_e - n_i), \tag{4}$$

which can be written in nondimensional form as

$$\frac{\partial n_i}{\partial t} + \frac{\partial(n_i v_{ix})}{\partial x} + \frac{\partial(n_i v_{iy})}{\partial y} + \frac{\partial(n_i v_{iz})}{\partial z} = 0, \tag{5}$$

$$\frac{dv_{ix}}{dt} = -\frac{\partial\phi}{\partial x} + bv_{iy} - \nu v_{ix}, \tag{6}$$

$$\frac{dv_{iy}}{dt} = -\frac{\partial \phi}{\partial y} - bv_{ix} - \nu v_{iy}, \tag{7}$$

$$\frac{dv_{iz}}{dt} = -\frac{\partial \phi}{\partial z} - \nu v_{iz}, \tag{8}$$

$$n_e^{2/3} = 1 + 2\phi + \frac{H^2}{\sqrt{n_e}} \nabla^2 \sqrt{n_e}, \tag{9}$$

$$\nabla^2 \phi = n_e - n_i. \tag{10}$$

Here, $d/dt \equiv \partial/\partial t + v_{ix}\partial/\partial x + v_{iy}\partial/\partial y + v_{iz}\partial/\partial z$ is the total derivative. In the equation $n_{e(i)}$ is the electron (ion) number density normalized by the equilibrium density n_0, $v_i \equiv (v_{ix}, v_{iy}, v_{iz})$ is the ion velocity normalized by the quantum IA speed $c_s = \sqrt{2k_B T_{Fe}/m_i}$, with k_B denoting the Boltzmann constant, m_i is the ion mass, $T_{Fe} = \hbar^2 (3\pi^2 n_0)^{2/3}/2k_B m_e$ is the electron Fermi temperature and $\hbar$ is the Planck's constant. Also, γ_i is the ion-neutral collision frequency and $\nu = \gamma_i/\omega_{pi}$ is the nondimensional collisional parameter. $H = \frac{\hbar \omega_{pe}}{2k_B T_{Fe}}$ is the quantum diffraction parameter denoting the ratio of the 'plasmon energy density' to the Fermi thermal energy in which $\omega_{pj} = n_0 e^2/\epsilon_0 m_j$ is the plasma oscillation frequency for the jth species particle. Moreover, ϕ is the electrostatic potential normalized by $2k_B T_{Fe}/e$. The space and time variables are respectively, normalized by c_s/ω_{pi} and the ion plasma period ω_{pi}^{-1}. The space–dependent parameter $b(\mathbf{r})$ is defined by $b(\mathbf{r}) = \frac{eB(\mathbf{r})}{m_i \omega_{pi}}$. We note that Eq. (9) is obtained after integrating once the momentum Eq. (3) for nonrelativistic degenerate electrons and using the boundary conditions : $\phi \to 0$, $n_e \to 1$ at infinity. In this equation, we have considered the pressure p_e follows the equation of state for degenerate electrons [45, 46]

$$p_e = \frac{1}{5} \frac{m_e V_{Fe}^2}{n_0^{2/3}} n_e^{5/3}, \tag{11}$$

where $V_{Fe} = \sqrt{2k_B T_{Fe}/m_e}$ is the electron Fermi thermal speed.

In order to describe the evolution of the nonlinear waves, we use the standard perturbation techniques. We define stretched variables,

$$\xi = \epsilon^{1/2}(l_x x + l_y y + l_z z - \lambda t), \quad \tau = \epsilon^{3/2} t, \tag{12}$$

where ϵ is a small parameter representing the strength of the wave amplitude, λ is the wave phase velocity and l_x, l_y, l_z are the direction cosines of the wave vector $\mathbf{k}$ having x, y, and z-components, i.e., $l_x^2 + l_y^2 + l_z^2 = 1$. We also assume that $\gamma_i = \epsilon^{3/2}\gamma$, where γ is of the order of unity or less. Such a consideration of smallness of γ_i can be found in the literature [47] and is valid in many experimental situations [48]. We now expand the dependent variables in a power series in terms of an expansion parameter ϵ,

$$\left.\begin{aligned}
n_i &= 1 + \epsilon n_i^{(1)} + \epsilon^2 n_i^{(2)} + \cdots, \\
n_e &= 1 + \epsilon n_e^{(1)} + \epsilon^2 n_e^{(2)} + \cdots, \\
v_{ix} &= \epsilon^{3/2} v_{ix}^{(1)} + \epsilon^2 v_{ix}^{(2)} + \cdots, \\
v_{iy} &= \epsilon^{3/2} v_{iy}^{(1)} + \epsilon^2 v_{iy}^{(2)} + \cdots, \\
v_{iz} &= \epsilon v_{iz}^{(1)} + \epsilon^2 v_{iz}^{(2)} + \cdots, \\
\phi &= \epsilon \phi^{(1)} + \epsilon^2 \phi^{(2)} + \cdots,
\end{aligned}\right\} \tag{13}$$

where ϵ is a smallness parameter measuring the strength of nonlinearity. Inserting the above stretching and expansions and comparing the lowest power of ϵ, we have

$$n_e^{(1)} = n_i^{(1)} = 3\phi^{(1)}, \quad v_{ix}^{(1)} = -\frac{l_y}{b} \frac{\partial \phi^{(1)}}{\partial \xi},$$

$$v_{iy}^{(1)} = \frac{l_x}{b} \frac{\partial \phi^{(1)}}{\partial \xi}, \quad v_{iz}^{(1)} = \frac{3\lambda}{l_z} \phi^{(1)},$$

together with

$$\lambda = \frac{|l_z|}{\sqrt{3}}, \tag{14}$$

where $l_z = (\mathbf{k}.\hat{z})/k = \cos\theta$ in which θ is the obliqueness angle between the wave propagation direction and the external magnetic field. The dynamical equations in the next higher order of ϵ are obtained as

$$\begin{aligned}
\lambda \frac{\partial n_i^{(2)}}{\partial \xi} - l_x \frac{\partial v_{ix}^{(2)}}{\partial \xi} - l_y \frac{\partial v_{iy}^{(2)}}{\partial \xi} - l_z \frac{\partial v_{iz}^{(2)}}{\partial \xi} = \frac{\partial n_i^{(1)}}{\partial \tau} \\
+ l_z \frac{\partial \left(n_i^{(1)} v_{iz}^{(1)}\right)}{\partial \xi}
\end{aligned} \tag{15}$$

$$v_{ix}^{(2)} = \frac{\lambda}{b} \frac{\partial v_{iy}^{(1)}}{\partial \xi} \tag{16}$$

$$v_{iy}^{(2)} = -\frac{\lambda}{b} \frac{\partial v_{ix}^{(1)}}{\partial \xi} \tag{17}$$

$$\begin{aligned}
\lambda \frac{\partial v_{iz}^{(2)}}{\partial \xi} - l_z \frac{\partial \phi^{(2)}}{\partial \xi} = \frac{\partial v_{iz}^{(1)}}{\partial \tau} + l_z v_{iz}^{(1)} \frac{\partial v_{iz}^{(1)}}{\partial \xi} \\
+ \gamma v_{iz}^{(1)}
\end{aligned} \tag{18}$$

$$\phi^{(2)} = \frac{1}{3} n_e^{(2)} - \frac{1}{18} n_e^{(1)^2} - \frac{H^2}{4} \frac{\partial^2 n_e^{(1)}}{\partial \xi^2} \tag{19}$$

$$\frac{\partial^2 \phi^{(1)}}{\partial \xi^2} = n_e^{(2)} - n_i^{(2)} \tag{20}$$

Finally, eliminating all the second-order quantities from Eqs. (15)–(20) together with relation to Eq. (14), we obtain the following equation :

$$\frac{\partial\phi^{(1)}}{\partial\tau}+4\lambda\phi^{(1)}\frac{\partial\phi^{(1)}}{\partial\xi}+\frac{\lambda}{6}\left(1-\frac{9H^2}{4}\right)\frac{\partial^3\phi^{(1)}}{\partial\xi^3}+\frac{\lambda}{6}(1-3\lambda^2)\frac{\partial}{\partial\xi}\left(\frac{1}{b}\frac{\partial}{\partial\xi}\left(\frac{1}{b}\frac{\partial\phi^{(1)}}{\partial\xi}\right)\right)+\frac{\gamma}{2}\phi^{(1)}=0 \tag{21}$$

In the above damped modified KdV Eq. (21), the fourth term represents the effects of the important parameter b related to the variable magnetic field, and the fifth term corresponds to the effects of ion-neutral collision. In the absence of magnetic field and collision, we may obtain the usual KdV equation for an IAW in quantum plasma. It should be mentioned that, in absence of magnetic field ($b = 0$), the corresponding term does not appear in the equation of motion of ions and therefore we are not concerned about the singularity in $b = 0$. The above Eq. (21) can be written as

$$\frac{\partial\phi^{(1)}}{\partial\tau}+A\phi^{(1)}\frac{\partial\phi^{(1)}}{\partial\xi}+B\frac{\partial^3\phi^{(1)}}{\partial\xi^3}+C(\xi)\frac{\partial^3\phi^{(1)}}{\partial\xi^3}+D(\xi)\frac{\partial^2\phi^{(1)}}{\partial\xi^2}+E(\xi)\frac{\partial\phi^{(1)}}{\partial\xi}+\frac{\gamma}{2}\phi^{(1)}=0 \tag{22}$$

where

$$\left.\begin{aligned}&A=4\lambda,\quad B=\frac{\lambda}{6}\left(1-\frac{9H^2}{4}\right),\quad C(\xi)=\frac{\lambda}{6b^2}(1-3\lambda^2),\\&D(\xi)=\frac{3\lambda}{6b^3}(1-3\lambda^2)\frac{\partial b}{\partial\xi},\\&E(\xi)=\frac{\lambda}{6}(1-3\lambda^2)\left\{\frac{3}{b^4}\left(\frac{\partial b}{\partial\xi}\right)^2-\frac{1}{b^3}\frac{\partial^2 b}{\partial\xi^2}\right\}.\end{aligned}\right\} \tag{23}$$

From the above Eq. (22), it is found that, in a uniform magnetic field, the fifth and sixth terms do not appear in the evolution equation. But, in a spatially varying magnetic field, the coefficient of dissipative terms is a spatial function through the nonuniform magnetic field encoded in $b(\xi)$. These two terms give rise to the generation of shocks and the frictional force term (due to the effects of ion-neutral collision) provides damping of the wave.

3 Results and discussion

In this section, we are interested in finding the solution of Eq. (22) with its full generality. In the presence of nonuniform magnetic field and collision, Eq. (22) is not an exactly integrable Hamiltonian system. Therefore, to investigate the effect of spatially varying magnetic field and collision on nonlinear low-frequency waves, we solve the nonlinear Eq. (22) numerically with the help of a MATHEMATICA-based finite difference scheme. As is well known, in the absence of magnetic field and collision, the nonlinear Eq. (22) is an exactly integrable Hamiltonian system and possesses a single soliton solution, $\phi^{(1)}(\xi,\tau)=\frac{3U}{A}sech^2\left[\frac{\xi-U\tau}{\sqrt{4B/U}}\right]$, where U is the soliton velocity. In order to examine the influence of spatially varying magnetic field on the dynamical properties of IAWs, we carry out such a numerical investigation of Eq. (22) for an arbitrary Gaussian shape of the magnetic field as $b=0.4(1+e^{-\xi^2})$.

For the time-dependent numerical solution, we use the single soliton solution as the initial wave form: $\phi^{(1)}(\xi,0)=\frac{3U}{A}sech^2[\sqrt{U/4B}\,\xi]$, $\xi\in[-L,L]$, where L is the spatial length. The boundary conditions are $\phi^{(1)}(\pm L,\tau)=\frac{3U}{A}sech^2[\sqrt{U/4B}\,L]$ and $\phi^{(1)}_\xi(-L,\tau)=0=\phi^{(1)}_\xi(L,\tau)$. To obtain ample results for the computation, we take $L = 100$ and $U = 0.4$. The development of this waveform at different times and for several values of parameters is displayed in Figures 1–4. We have analyzed the influence of different spatially varying magnetic field perturbations in absence of collision on the time evolution of nonlinear structures in Figure 1. It is found that as the time progresses, the pulse amplitude decreases. One also sees the generation of oscillatory tail (dispersive shock) occur behind the advancing soliton. As the magnetic field strength decreases, the number of radiating oscillatory structures increases and the wave amplitude reduces prominently in its magnitude. This figure clearly indicates that smaller magnetic field creates larger dissipation effects. This is expected as the dissipation coefficient is proportional to $1/b^3$ according to Eq. (23). Figure 2(a) and (b) show the development of wave profile in absence and presence of collision, respectively. This clearly demonstrates the generation of oscillatory tail at different time behind the soliton. In the absence of collisional effect [Figure 2(a)], the shock is more dispersive than that of the presence of collision [Figure 2(b)]. Also, it is seen that in the presence of collision the decrease in soliton amplitude become more pronounced. This clearly demonstrates that the presence of collision enhances the dissipative effect. Figure 3 depicts the developments of solitary waves into shocks for different values of quantum diffraction parameter H. It is observed that the solitary pulse becomes smaller in amplitude and narrower in width for higher values of H. It means that quantum effects compress the soliton. This may be attributed to the fact that the increase in quantum diffraction parameter H reduces dispersion in the system, and consequently the amplitude is decreased. One of the important results of the soliton perturbation is the formation of oscillatory tail (a wave packet of small amplitude following behind a soliton). It may be mentioned that for larger values of H the amplitude of this radiative shock structures increases. Finally, we have examined the influence of collisional parameter γ on the basic features of nonlinear wave profile in Figure 4. It is evident that due to collisionality, the wave is damped and its amplitude reduces as it propagates. Also, we note that as the value of collisional parameter increases, the number of oscillatory tail reduces. This is because that shock is created due to dissipation, and increasing collision frequency is commensurate to enhancing the dissipation in the system. Therefore, collisional effect significantly modifies the nonlinear characteristics of IAW profiles.

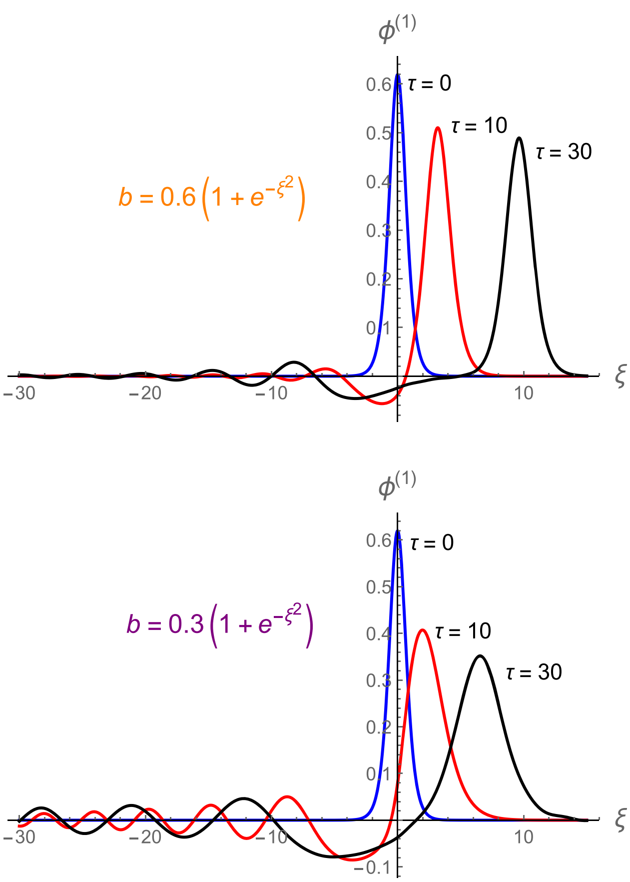

Figure 1: Solution of Eq. (22) in different time τ in the absence of collision for different varying magnetic fields, where $H = 0.1$, $\theta = 10°$ and $U = 0.4$.

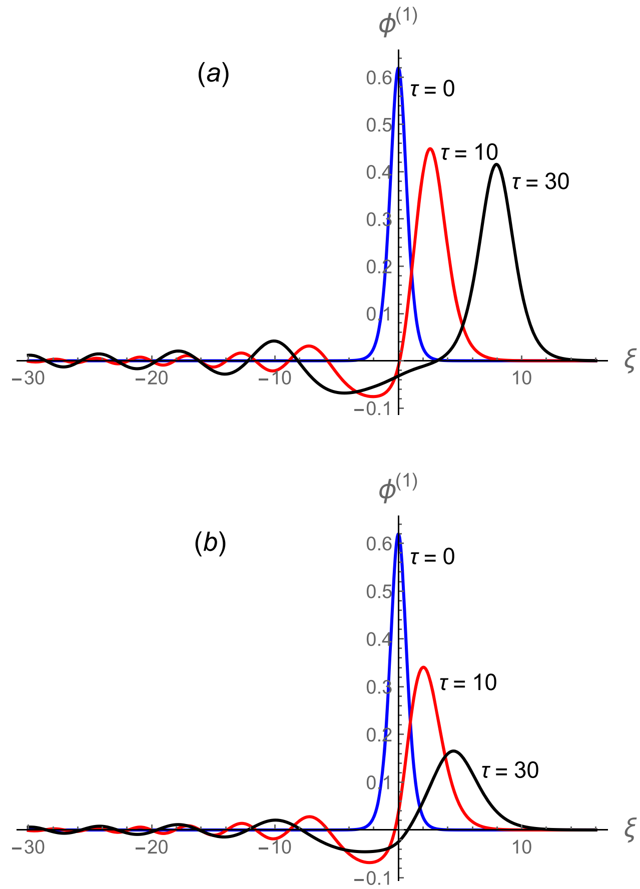

Figure 2: Development of soliton structures at different time (a) in absence and (b) presence of collision ($\gamma = 0.05$), where $b = 0.4\,(1 + e^{-\xi^2})$ and other parameters are same as in Figure 1.

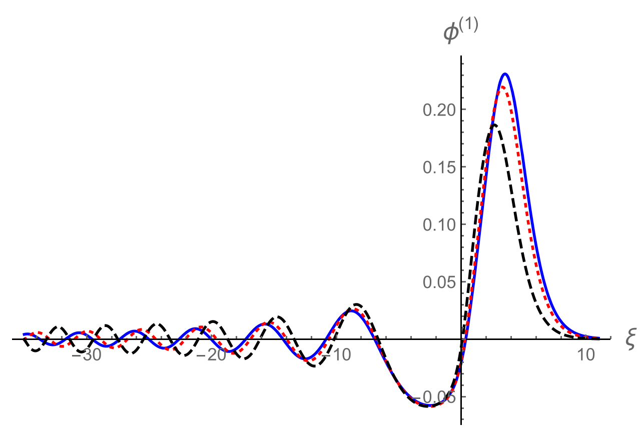

Figure 3: The effect of quantum diffraction parameter H on the nonlinear excitation structures: solid line for $H = 0.1$, dotted line for $H = 0.3$, dashed line for $H = 0.5$, where $\gamma = 0.05$, $b = 0.4\,(1 + e^{-\xi^2})$ and other parameters are same as in Figure 1.

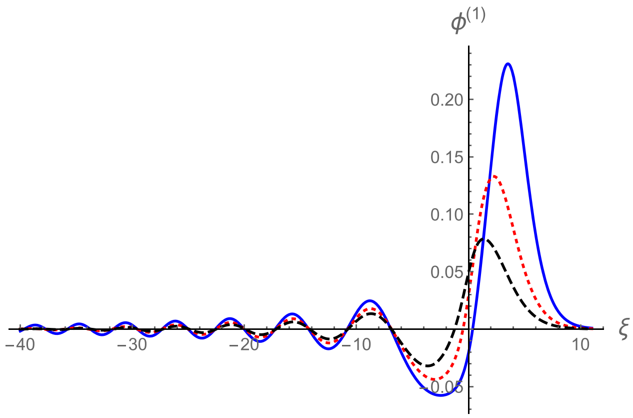

Figure 4: The effect of collisional parameter γ on the nonlinear excitation structures: solid line for $\gamma = 0.05$, dotted line for $\gamma = 0.1$, dashed line for $\gamma = 0.15$, where $b = 0.4\,(1 + e^{-\xi^2})$ and other parameters are same as in Figure 1.

4 Conclusion

We have investigated the propagation characteristics of low-frequency nonlinear electrostatic waves in a dissipative quantum plasma. The dissipation aries due to the spatially varying magnetic field and ion-neutral collision. The dynamics of the nonlinear wave is found to be governed by a nonlinear partial differential equation in the form of modified KdVB equation, by adopting the reductive perturbation technique in the quantum plasma medium under consideration. In our analysis, we have focused particularly on the time-dependent amplitude and generation of oscillatory tail in order to trace the effect of dissipation. We have shown that the presence of nonuniform magnetic field is responsible for the Burgers' term. This brings the physics of shock wave and plays a significant role in propagation dynamics. Till now, we have not yet envisaged any experimental observations. We hope that the next generation intense laser plasma laboratory experiments will furnish the experimental investigation of these types of nonlinear structures (oscillatory tail behind the soliton) in quantum plasmas.

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