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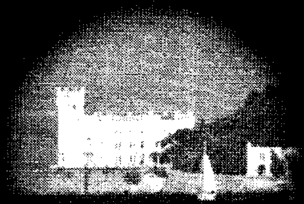

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Available at: http://www.ictp.trieste.it/~pub_off IC/99/61

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EFFECTS OF DUST GRAIN CHARGE FLUCTUATION ON OBLIQUELY PROPAGATING DUST-ACOUSTIC POTENTIAL IN MAGNETIZED DUSTY PLASMAS

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Abstract

Effects of dust grain charge fluctuation, obliqueness and external magnetic field on finite amplitude dust-acoustic solitary potential in a magnetized dusty plasma, consisting of electrons, ions and charge fluctuating dust grains, have been investigated by the reductive perturbation method. It has been shown that such a magnetized dusty plasma system may support dust-acoustic solitary potential on a very slow time scale involving the motion of dust grains, whose charge is self-consistently determined by local electron and ion currents. The effects of dust grain charge fluctuation, external magnetic field and obliqueness are found to modify the properties of this dust-acoustic solitary potential significantly. The implications of these results to some space and astrophysical dusty plasma systems, especially to planetary ring-systems and cometary tails, are briefly mentioned.

MIRAMARE – TRIESTE

May 1999

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1. Introduction

Recently, there has been a great deal of interest in understanding different types of collective processes in dusty plasmas (plasmas with extremely massive and negatively charged dust grains), because of its vital role in the study of astrophysical and space environments, such as, asteroid zones, planetary atmospheres, interstellar media, circumstellar disks, dark molecular clouds, cometary tails, nebulae, earth's environment, etc. [1-7]. These dust grains are invariably immersed in the ambient plasma and radiative background. The interaction of these dust grains with the other plasma particles (viz. electrons and ions) is due to the charge carried by them. The dust grain is charged by a number of competing processes, depending upon the local conditions, such as, photoelectric emission stimulated by the ultra-violet radiation, collisional charging by electrons and ions, disruption and secondary emission due to the Maxwellian stress, etc. [8-12].

It has been found that the presence of static charged dust grains modifies the existing plasma wave spectra [13-20]. Bliokh and Yaroshenko [13] studied electrostatic waves in dusty plasmas and applied their results in interpreting spoke-like structures in Saturn's rings (revealed by Voyager space mission [21]). Angelis *et al.* [14] investigated the propagation of ion acoustic waves in a dusty plasma, in which a spatial inhomogeneity is created by a distribution of immobile dust particles [22]. They [14] applied their results in interpreting the low frequency noise enhancement observed by the *Vega and Giotto* space probes in the dusty regions of Halley's comet [23].

On the other hand, it has been shown both theoretically and experimentally that the dust charge dynamics introduces new eigenmodes [24-33]. The dust-acoustic mode [24-28], where dust particle mass provides the inertia and the pressures of electrons and ions give rise to the restoring force, is one of them. Rao *et al.* [24] have first reported theoretically the existence of extremely low phase velocity (in comparison with the electron and ion thermal velocities) dust-acoustic waves in an unmagnetized dusty plasma. The laboratory experiments of Barkan *et al.* [27] and D'Angelo [28] have conclusively verified this theoretical prediction of Rao *et al.* [24] and reported some nonlinear features of the dust-acoustic waves.

Recently, motivated by the experimental observations [27,28] of these low phase velocity dust-acoustic waves, a number of investigations [34-38] have been made to study nonlinear properties of these novel dust-acoustic waves. Most of these works have considered an unmagnetized dusty plasma system containing dust grains of constant negative charge. The assumption of the dust grains having constant negative charge in such a plasma system virtually represents a plasma with heavy species of negative ions. However, the charge of the dust grain varies according to the local plasma currents flowing into the surface of these grains. Thus, the grain charge is a variable which has to be determined self-consistently by its charging currents. On the other hand, it is well known that the effects of external magnetic field and obliqueness, which have also not been considered in these earlier investigations [34-38], drastically modify the properties of the electrostatic solitary structures [39-41]. Thus, to examine the effects of the dust grain charge fluctuation, external magnetic field and obliqueness, in the present work we have studied the obliquely propagating dust-acoustic solitary waves in a magnetized dusty plasma system consisting of electrons, ions and dust grains whose charge is self consistently determined by the local electron and ion currents flowing to the surface of these dust grains.

The paper is organized as follows. The basic equations governing the plasma system under consideration is presented in section 2. The Korteweg de Vries (K-dV) equation is derived by employing the reductive perturbation method in section 3. The solitary wave solution of this K-dV equation is obtained and the properties of these electrostatic solitary potential are discussed in section 4. Finally, a brief discussion is given in section 5.

2. Governing equations

We consider a three component magnetized dusty plasma system consisting of electrons, ions and cold, inertial, charge fluctuating dust grains whose charge varies according to the local plasma currents flowing into the surface of these grains. This plasma system is assumed to be immersed in an external static magnetic field ($\mathbf{B}_0 \parallel \hat{\mathbf{z}}$). On the extremely slow dust time scale, the electrons and ions are in local thermodynamic equilibrium and their number densities, n_e and n_i, obey the Boltzmann distribution [3, 24-26]:

$$\left.\begin{array}{l} n_e = n_{eo} \exp\left(\frac{e\varphi}{k_B T_e}\right), \\ n_i = n_{io} \exp\left(-\frac{e\varphi}{k_B T_i}\right), \end{array}\right\} \tag{1}$$

where n_{e0} (n_{i0}) and T_e (T_i) are the equilibrium number density and temperature of electrons (ions), respectively; φ is the electrostatic wave potential; e is the magnitude of the electron charge; k_B is the Boltzmann constant. The dust particles are all assumed to be spherical with the same radius a_d and surface charge q_d. If the dust particles are charged by plasma currents due to the colliding electrons and ions, we have for the electron (ion) current I_e (I_i) [3]:

$$\left.\begin{array}{l} I_e(q_d,\varphi) = -\pi a_d^2 e(\frac{8k_BT_e}{\pi m_e})^{1/2} n_e \exp\left(\frac{eq_d}{a_d k_B T_e}\right), \\ I_i(q_d,\varphi) = \pi a_d^2 e(\frac{8k_BT_i}{\pi m_i})^{1/2} n_i \exp\left(-\frac{eq_d}{a_d k_B T_i}\right), \end{array}\right\} \tag{2}$$

where m_e (m_i) is the mass of an electron (ion). The dynamics of low phase velocity (lying between the ion and dust thermal velocities, viz., $v_{td} << v_p << v_{ti}$) dust-acoustic oscillations is governed by [24-26]

$$\frac{\partial n_d}{\partial t} + \nabla\cdot(n_d \mathbf{u_d}) = 0, \tag{3}$$

$$\frac{\partial \mathbf{u_d}}{\partial t} + (\mathbf{u_d}\cdot\nabla)\mathbf{u_d} = -\frac{q_d}{m_d}\nabla\varphi + \frac{q_d}{m_d c}(\mathbf{u}_d \times \mathbf{B}_0), \tag{4}$$

$$\frac{\partial q_d}{\partial t} + (\mathbf{u_d}\cdot\nabla)q_d = I_e(q_d,\varphi) + I_i(q_d,\varphi), \tag{5}$$

$$\nabla^2\varphi = 4\pi[e(n_e - n_i) - q_d n_d], \tag{6}$$

where n_d is the dust particle number density; $\mathbf{u}_d$ is the dust fluid velocity; m_d is dust particle mass; c is the speed of light in vacuum.

3. Derivation of the K-dV equation

To examine the nature of the dust-acoustic solitary potential in our dusty plasma model, we construct a weakly nonlinear theory of the dust-acoustic waves with small but finite amplitude which leads to the scaling of the independent variables through the stretched coordinates [41,42]

$$\left.\begin{array}{l} \xi = \epsilon^{1/2}(l_x x + l_y y + l_z z - v_0 t), \\ \tau = \epsilon^{3/2} t, \end{array}\right\} \tag{7}$$

where ϵ is a small parameter measuring the weakness of the dispersion, v_0 is the wave phase velocity; l_x, l_y and l_z are the directional cosines of the wave vector $\mathbf{k}$ along the x-, y- and z-axes, respectively, so that $l_x^2 + l_y^2 + l_z^2 = 1$. We can expand the perturbed quantities n_d, q_d, φ and u_{dz} about their equilibrium values (equilibrium values of n_d and q_d are n_{d0} and q_{d0}, respectively) in power of ϵ by following Refs. 41 and 42. To obtain x- and y-components of dust electric field and polarization drifts, we can expand the perturbed quantities $u_{dx,y}$ by following a standard technique [41] where the terms of $\epsilon^{3/2}$ are included. Thus, we can expand n_d, q_d, φ and $u_{dx,y,z}$ as [40,41]

$$\left.\begin{array}{l} n_d = n_{d0} + \epsilon n_d^{(1)} + \epsilon^2 n_d^{(2)} + \cdots\cdots, \\ q_d = q_{d0} + \epsilon q_d^{(1)} + \epsilon^2 q_d^{(2)} + \cdots\cdots, \\ \varphi = 0 + \epsilon\varphi^{(1)} + \epsilon^2\varphi^{(2)} + \cdots\cdots, \\ u_{dz} = 0 + \epsilon u_{dz}^{(1)} + \epsilon^2 u_{dz}^{(2)} + \cdots\cdots, \\ u_{dx,y} = 0 + \epsilon^{3/2} u_{dx,y}^{(1)} + \epsilon^2 u_{dx,y}^{(2)} + \cdots\cdots. \end{array}\right\} \tag{8}$$

Now, using (7) and (8) in (5) and (6) [with (1) and (2)], we obtain zeroth order equations which give

$$\left.\begin{array}{l} I_e(\varphi = 0, q_d = q_{d0}) + I_e(\varphi = 0, q_d = q_{d0}) = 0, \\ en_{e0} - en_{i0} - q_{d0} n_{0d} = 0. \end{array}\right\} \tag{9}$$

If we set $q_{d0} = -Z_d e$ (where Z_d is the number of electrons residing on the dust grain at equilibrium), at equilibrium we have $n_{io} = Z_d n_{d0} + n_{e0}$. This is used by a number of authors, for example, Rao *et al.* [24], Mamun *et al.* [34,35], Kotsarenko *et al.* [37], Roychoudhury and Chatterjee [38], etc., who have neglected the dust grain charge fluctuation. Again, using (7) and (8) in (3) $-$ (6) one can obtain the first order perturbation equations which, after simplification, yield

$$\left.\begin{array}{l} u_{dz}^{(1)} = \frac{l_z}{v_0}(\frac{q_{d0}}{m_d})\varphi^{(1)}, \\ n_d^{(1)} = n_{d0}(\frac{l_z}{v_0})^2(\frac{q_{d0}}{m_d})\varphi^{(1)}, \\ q_d^{(1)} = -a_d\varphi^{(1)}, \\ v_0 = l_z\frac{C_0}{\sqrt{(1+\mu)}}, \end{array}\right\} \tag{10}$$

where $\mu = 4\pi a_d \lambda_{D0}^2 n_{d0}$ is a dimensionless parameter [10,25] which estimates whether the background dust cloud is tenuous or dense; $\lambda_{D0} = [\sum_{s=e,i}(4\pi n_{s0}e^2)/k_B T_s]^{-1/2}$; $C_0 = \omega_{p0}\lambda_{D0}$; $\omega_{p0} = (4\pi n_{d0} q_{d0}^2/m_d)^{1/2}$. It should be mentioned here that the last of Eqs. (10) represents the linear dispersion relation for the undamped dust-acoustic mode where the effects of the dust grain charge fluctuation and obliqueness are included. However, one can obtain a linear dispersion relation for the damped dust acoustic mode [25] by taking Fourier images of (3) - (6), instead of using the stretched coordinates we used. It is obvious from our dispersion relation that the effect of the charge fluctuation of the dust grain decreases the wave phase velocity of the dust-acoustic mode. If we put $\mu = 0$ and $l_z = 1$, this dispersion relation corresponds to that obtained by Rao *et al.* [24] and by some others [34-38]. We can write the first order x- and y-components of the momentum equation as

$$\left.\begin{aligned} u_{dy}^{(1)} &= c(\tfrac{l_z}{B_0})\tfrac{\partial\varphi^{(1)}}{\partial\xi}, \\ u_{dx}^{(1)} &= -c(\tfrac{l_y}{B_0})\tfrac{\partial\varphi^{(1)}}{\partial\xi}. \end{aligned}\right\} \tag{11}$$

These, respectively, represent the y- and x-components of the electric field drift. These equations are also satisfied by the second order continuity equation.

Again, using (7) and (8) in (5), and eliminating $u_{dx,y}^{(1)}$, we obtain the next higher order x- and y-components of the momentum equation as

$$\left.\begin{aligned} u_{dy}^{(2)} &= \tfrac{l_y v_0}{\omega_{c0}^2}(\tfrac{q_{d0}}{m_d})\tfrac{\partial^2\varphi^{(1)}}{\partial\xi^2}, \\ u_{dx}^{(2)} &= \tfrac{l_x v_0}{\omega_{c0}^2}(\tfrac{q_{d0}}{m_d})\tfrac{\partial^2\varphi^{(1)}}{\partial\xi^2}, \end{aligned}\right\} \tag{12}$$

where $\omega_{c0} = |q_{d0}|B_0/m_d c$. These equations, respectively, denote the y- and x-components of the dust polarization drift. We can also express the second order perturbation equations for (5) and (6) as

$$\left.\begin{aligned} q_d^{(2)} &= -a_d\varphi^{(2)} + \tfrac{1}{2}(1-\tfrac{T_i}{T_e})\tfrac{ea_d}{k_BT_i}[\varphi^{(1)}]^2 + (1-\tfrac{T_i}{T_e})\tfrac{e}{k_BT_i}q_d^{(1)}\varphi^{(1)} + \tfrac{1}{2}(1-\tfrac{T_i}{T_e})\tfrac{e}{a_dk_BT_i}[q_d^{(1)}]^2, \\ \tfrac{\partial^2\varphi^{(1)}}{\partial\xi^2} &= \tfrac{1}{\lambda_{D0}^2}\varphi^{(2)} - 4\pi q_{d0}n_d^{(2)} - 4\pi n_{d0}q_d^{(2)} - 4\pi q_d^{(1)}n_d^{(1)} - \tfrac{1}{2}\tfrac{e\alpha}{k_BT_i\lambda_{Di}^2}[\varphi^{(1)}]^2, \end{aligned}\right\} \tag{13}$$

where $\alpha = 1 - (n_{e0}/n_{i0})(T_i^2/T_e^2)$ and $\lambda_{Di} = (k_BT_i/4\pi n_{i0}e^2)^{1/2}$. Similarly, following the same procedure one can obtain the next higher order continuity equation and z-component of the momentum equation as

$$\left.\begin{aligned} &\tfrac{\partial n_d^{(1)}}{\partial\tau} - v_0\tfrac{\partial n_d^{(2)}}{\partial\xi} + n_{d0}\tfrac{\partial}{\partial\xi}[l_xu_{dx}^{(2)} + l_yu_{dy}^{(2)} + l_zu_{dz}^{(2)}] + l_z\tfrac{\partial}{\partial\xi}[n_d^{(1)}u_{dz}^{(1)}] = 0, \\ &\tfrac{\partial u_{dz}^{(1)}}{\partial\tau} - v_0\tfrac{\partial u_{dz}^{(2)}}{\partial\xi} + l_zu_{dz}^{(1)}\tfrac{\partial u_{dz}^{(1)}}{\partial\xi} + l_z(\tfrac{q_{d0}}{m_d})\tfrac{\partial\varphi^{(2)}}{\partial\xi} + l_z(\tfrac{1}{m_d})q_d^{(1)}\tfrac{\partial\varphi^{(1)}}{\partial\xi} = 0. \end{aligned}\right\} \tag{14}$$

Now, using (9) $-$ (14), one can eliminate $n_d^{(2)}$, $q_d^{(2)}$, $u_{dz}^{(2)}$ and $\varphi^{(2)}$. The elimination of these second order variables reduces to a nonlinear equation of the form

$$\frac{\partial\varphi^{(1)}}{\partial\tau} + a\varphi^{(1)}\frac{\partial\varphi^{(1)}}{\partial\xi} + b\frac{\partial^3\varphi^{(1)}}{\partial\xi^3} = 0. \tag{15}$$

This is known as K-dV equation with nonlinear coefficient a and dispersion coefficient b which, in our case, are given by

$$\left.\begin{aligned} a &= \tfrac{3}{2}(\tfrac{eC_0}{k_BT_i})(\tfrac{q_d0}{Z_de})(\tfrac{\sigma l_z}{\beta\sqrt{1+\mu}})[1 + \tfrac{1}{3}(\tfrac{Z_de}{q_{d0}})\tfrac{\alpha\beta}{\sigma^2(1+\mu)}], \\ b &= \tfrac{1}{2}\tfrac{l_z}{(1+\mu)^{3/2}}(C_0\lambda_{D0}^2)[1 + \tfrac{\omega_{p0}^2}{\omega_{c0}^2}(1-l_z^2)], \end{aligned}\right\} \tag{16}$$

where $\beta = (1 - n_{e0}/n_{i0})$ and $\sigma = 1 + (n_{e0}/n_{i0})(T_i/T_e)$.

4. Solitary wave solution of the K-dV equation

The steady state solution of this K-dV equation is obtained by transforming the independent variables ξ and τ to $\eta = \xi - U_0\tau$ and $\tau = \tau$, where U_0 is a constant velocity, and imposing the appropriate boundary conditions, viz., $\varphi \to 0$, $\frac{d\varphi^{(1)}}{d\eta} \to 0$, $\frac{d^2\varphi^{(1)}}{d\eta^2} \to 0$ at $\eta \to \pm\infty$. Thus, one can express the steady state solution of this K-dV equation as

$$\varphi^{(1)} = \varphi_m^{(1)} \mathrm{sech}^2[(\xi - U_0\tau)/\Delta], \tag{17}$$

where the amplitude $\varphi_m^{(1)}$ and the width Δ are given by

$$\left.\begin{array}{l} \varphi_m^{(1)} = 3U_0/a, \\ \Delta = \sqrt{4b/U_0}. \end{array}\right\} \tag{18}$$

We now normalize the amplitude $\varphi_m^{(1)}$ by $k_B T_i/e$, i.e., $e\varphi_m^{(1)}/k_B T_i = \psi$, Δ by λ_{D0}, i.e., $\Delta/\lambda_{D0} = \delta$, velocity U_0 by C_0, i.e., $U_0/C_0 = u_0$ and ω_{c0} by ω_{p0}, i.e., $\omega_{c0}/\omega_{p0} = \omega_c$. These normalizations allow us to express the normalized amplitude and width of these solitary waves in simple form as

$$\left.\begin{array}{l} \psi = -2(\frac{u_0}{l_z})\frac{\beta}{\sigma}(\sqrt{1+\mu})\left(1 - \frac{1}{3}\frac{\alpha\beta}{\sigma^2}\frac{1}{1+\mu}\right)^{-1}, \\ \delta = \sqrt{2}\frac{l_z}{u_0(1+\mu)^{3/4}}\left(1 + \frac{1-l_z^2}{\omega_c^2}\right), \end{array}\right\} \tag{19}$$

where $q_d = -Z_d e$ and $Z_d n_{d0} + n_{e0} = n_{i0}$ are used.

It is obvious that for any dusty plasma system $0 < \alpha \leq 1$, $0 < \beta \leq 1$, and $1 < \sigma \leq 2$. It can also be estimated by available data [3-12] that for tenuous dust cloud $\mu << 1$ and for a dense dusty plasma $\mu \geq 1$ or even $\mu >> 1$. Thus, for any possible values of α, β, σ and μ, the amplitude ψ is negative, i.e., $\psi < 0$. This means that the dusty plasma system under consideration may support solitary waves with negative potential but not with positive potential. It is also seen from Eqs. (19) that for $\mu << 1$ the dust grain charge fluctuation does not have significant effect on formation or properties of these solitary structures but for $\mu >> 1$, this effect significantly modifies the properties of the solitary structures. It is shown that for $\mu >> 1$, $\psi \propto \sqrt{\mu}$ and $\delta \propto \mu^{-3/4}$.

It is obvious that the amplitude of these solitary waves is inversely proportional to l_z ($l_z = \cos\theta$, where θ is the angle between the directions of the wave propagation vector $\mathbf{k}$ and the external magnetic field $\mathbf{B}_0$). The width of these solitary waves is a nonlinear function of l_z an ω_c. The adjoining figure shows how the width (δ) of these solitary waves changes with the obliqueness (θ) and with the magnitude of the external magnetic field (ω_c). It is observed that the width (δ) increases with θ for its lower range (i.e. from 0° to ~ 45°), but decreases for its higher range (i.e. from ~ 45° to 90°). Though the variation of δ with θ has been shown for any value of θ in between 0° and 90°, our perturbation method (which is only valid for small but finite amplitude limit) is not valid for large θ which makes the wave amplitude large enough to beak the validity of the reductive perturbation method used.

It is shown from Eqs. (19) that if we have neglected the effects of grain charge fluctuation, external magnetic field, and obliqueness (i.e., we set $\mu = 0$ and $l_z = 1$), the expressions for the peak amplitude and width of these solitary waves corresponds to those obtained by a number of recent published works [24,34-38].

5. Discussion

A three component magnetized dusty plasma system, consisting of Boltzmann distributed electrons and ions, and charge fluctuating dust grains, has been considered and the properties of finite amplitude dust-acoustic solitary potential, which has been found to exist in such a dusty plasma system, have been investigated by the reductive perturbation method. The results, which have been found in this investigation, may be pointed out as follows:

i) It has been shown that for the study of electrostatic waves, the variation of the dust grain charge must be taken into account and that the phase velocity of the waves will be affected by the self-consistent dust grain charge variation. It is seen that this dust grain charge variation reduces the phase velocity of the dust-acoustic waves.

ii) We have shown that on a very slow time scale of the dust motion and with self consistent grain charge variation, a magnetized dusty plasma system may support obliquely propagating stable dust-acoustic solitary waves with negative potential only; the system does not support solitary waves with positive potential.

iii) It is also seen that for a tenuous dust cloud, where $\mu << 1$, the dust grain charge variation does not have significant effect on formation or properties of these electrostatic solitary structures but for a dense dusty plasma, where $\mu >> 1$, this effect significantly modifies the properties of these solitary structures. It is shown that for $\mu >> 1$, as we increase μ, the peak amplitude increases whereas the width decreases. This means that the dust grain charge variation makes solitary potential more spiky.

iv) It is obvious that the amplitude of these solitary waves is inversely proportional to l_z ($l_z = \cos\theta$). It is observed that the width (δ) increases with θ for its lower range (i.e. from $0°$ to $\sim 45°$), but decreases for its higher range (i.e. from $\sim 45°$ to $90°$). It should be pointed out that for large angles the assumption that the waves are electrostatic is no longer a valid one, and we should look for fully electromagnetic structures.

v) It is seen that the magnitude of the external magnetic field has no effect on the amplitude of the solitary potential. However, it does have an effect on the width of the solitary potential. It is shown that as we increase the magnitude of the magnetic field, the width of the solitary potential decreases, i.e., the external magnetic field also makes the solitary potential more spiky.

We have analyzed properties of the electrostatic solitary wave potential in a magnetized dusty plasma by the reductive perturbation method which is valid for small but finite amplitude limit. Since in many astrophysical situations there may exist extremely large amplitude solitary waves, we propose to develop a more exact theory for the study of obliquely propagating arbitrary-amplitude solitary waves in such a magnetized dusty plasma system. However, our present analysis should be useful for understanding different nonlinear features of localized electrostatic disturbances in a number of astrophysical dusty plasma systems, such as, planetary ring systems (viz. Saturn's rings [3,13,21]), cometary environment (viz. Halley's comet [14,23]), interstellar medium [3], etc., where charge fluctuating dust grains and Boltzmann distributed electrons and ions are the major plasma species.

To conclude it may be added that the time evolution and stability analysis of obliquely propagating arbitrary amplitude solitary structures, are also problems of great importance but beyond the scope of the present work.

Acknowledgements

AAM would like to express his gratitude to Prof. R. A. Cairns, Prof. P. K. Shukla, Prof. M. Salimullah, and Dr. Y. Hayashi for their stimulating influence and helpful discussions during the course of this work. He would also like to acknowledge the duty leave granted by the authority of Jahangirnagar University and the financial support of the Abdus Salam ICTP. This work was done with the framework of the Asscociateship Scheme of the Abdus Salam International Centre for Theoretical Physics, Trieste, Italy. Financial support from the Swedish International Develepment Cooperation Agency is acknowledged.

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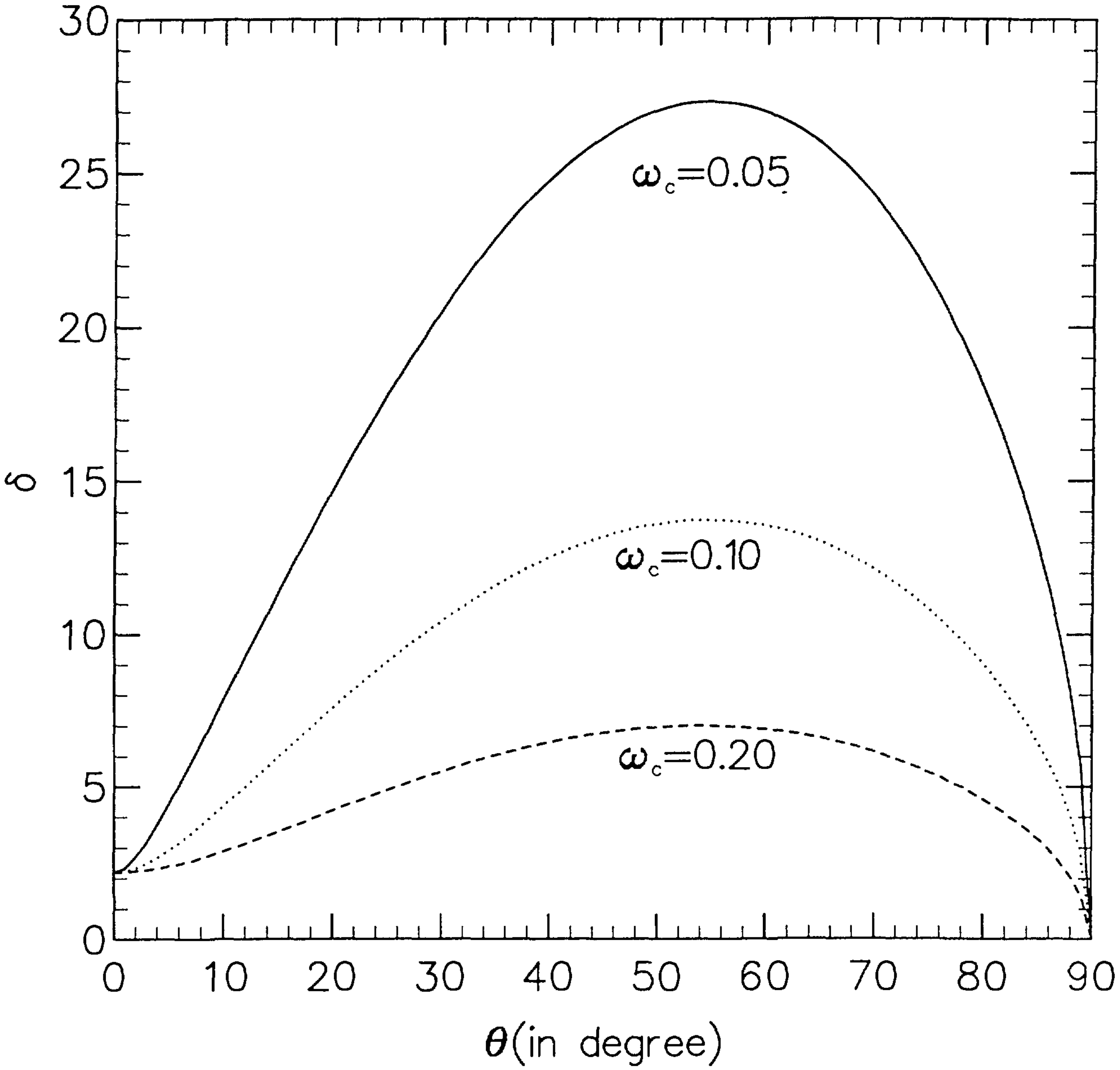

Fig.1

Variation of the width (δ) with the propagation angle (θ) for $\mu = 0.1$, $u_0 = 1.0$, $\omega_c = 0.05$ (solid curve), $\omega_c = 0.1$ (dashed curve), and $\omega_c = 0.2$ (dotted curve).