Propagation of dust-acoustic waves in a bounded dusty plasma

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The linear and nonlinear propagation of dust-acoustic waves in a dusty plasma bounded in finite geometry has been theoretically investigated. It is found that the finite geometry of the bounded plasma makes a significant contribution to the instability of the wave. Moreover, the pseudopotential has a positive and inverted profile that prevents the trapping of particles and does not favor solitary waves in bounded dusty plasma.

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I. INTRODUCTION

Studies on the propagation of waves in a dusty plasma have been found to be very important in space and astrophysical situations, fusion devices, industry, and modern technology [1–6]. A dusty plasma is a three component plasma with electrons, ions, and highly charged massive dust grains which influence the collective properties of the plasma. The size of the grains is typically in the micrometer range. The dust grains are generally negatively charged, $\sim 10^3 e - 10^5 e$, by various charging processes. The charge of the dust particles are not always fixed, but depends on the plasma properties, e.g., photoemission, etc. [7]. Variable charges introduce interesting characteristics in the dynamics of wave propagation in the plasma [8]. In the last few years, various authors investigated both linear and nonlinear wave propagation in plasma [9–14]. The authors of Ref. [15] were the first to investigate theoretically the collective motion of negatively charged dust particles in the background of hot electrons and ions in the plasma. They discussed a type of sound wave called the dust-acoustic wave of very low frequency, driven by plasma pressure and inertia. Subsequently, Refs. [16–22] studied theoretically the propagation of waves through a dusty plasma, and discussed different aspects of the dust acoustic mode under various physical situations. In the context of an ion acoustic solitary wave in a dusty plasma, the variable charges of the dust particles play important roles in the existence of solitons in the plasma. The authors of Ref. [23] showed that, for small grains, ion acoustic solitons may exist, but that an increase in the grain charge destroys them. Recently, it was shown [24] that the streaming motion and attachment coefficients of ions and electrons have significant effects on the formation of ion acoustic solitary wave in a dusty plasma.

It is worthwhile to mention that many authors studied the propagation of waves in a bounded plasma system [25–27], but no study has yet been done for the bounded dusty plasma system. Studies of dust acoustic waves in a bounded plasma system are very much important for its relevance in the context of laboratory experiments and industrial application. Our intention in the present paper is to investigate theoretically both linear and nonlinear wave propagation in a dusty plasma bounded in a finite geometry, together with a variable charge fluctuation of dust grains. From the linear dispersion relation it has been observed that the stability of a dust acoustic wave depends on the finite geometry of the bounded plasma system. Analyzing the basic equations with the help of the Sagdeev potential, it is found that the structure of the potential is completely different; this does not support the trapping of particles, and so a solitary wave cannot be formed in a bounded dusty plasma system.

II. BASIC EQUATIONS AND LINEAR DISPERSION RELATION

We consider a homogeneous, collisionless, and unmagnetized dusty plasma consisting of electrons, ions, and charged dust particles. The plasma is bounded in a finite geometry. We assume that a hydrodynamic description is possible for a plasma system where electrons and ions are in the background and are thermalized. The density distributions for the electrons and ions can be written as

$$n_e = n_{eo} \exp(e\phi/k_B T_e) \tag{1}$$

and

$$n_i = n_{io} \exp(-Z_i e \phi/k_B T_i), \tag{2}$$

where n_{eo} and n_{io} are the average equilibrium densities of electrons and ions, respectively.

The equations which govern the dynamical behavior of dust particles which are heavier than both electrons and ions are given by the cold fluid equations

$$(\partial/\partial t + \mathbf{v}_D \cdot \boldsymbol{\nabla})\mathbf{v}_D = -(Q_D/m_D)\boldsymbol{\nabla}\phi, \tag{3}$$

$$\partial n_D/\partial t + \boldsymbol{\nabla}\cdot(n_D \mathbf{v}_D) = 0, \tag{4}$$

where $\mathbf{v}_D$, n_D, m_D, and Q_D are the velocity, density, mass, and charge of the dust particles respectively.

The required Poisson equation is

$$\boldsymbol{\nabla}^2\phi = 4\pi e(n_e - Z_i n_i - n_D Q_D/e). \tag{5}$$

Moreover, the fluctuation of dust charges due to interactions of the dust with the electrons and ions is given by [8]

$$(\partial/\partial t + \mathbf{v}_D \cdot \mathbf{V})Q_D = I_e + I_i, \tag{6}$$

where I_e and I_i represent the electron current and the ion current:

$$I_e=-\pi a^2 e(8k_BT_e/\pi m_e)^{1/2} n_e \exp[(e/k_BT_e)(\phi_f-V_p)], \tag{7}$$

$$I_i=\pi a^2 e(8k_BT_i/\pi m_i)^{1/2} n_i Z_i\{1-(eZ_i/k_BT_i)(\phi_f-V_p)\}. \tag{8}$$

Here a is the radius of the dust particles. T_e (T_i), m_e (m_i), and n_e (n_i) are the temperature, mass, and density of electrons (ions). (ϕ_f-V_p) is the difference between the potentials of dust-grain surface and ambient plasma. It should be mentioned here that expressions for the electron and ion currents are useful only when the dust grains are stationary.

Now, for the study of wave propagation in a bounded plasma system, we write Eqs. (1)–(6) in normalized form as given below:

$$N_e=\alpha_{ne}\exp(\psi), \tag{9}$$

$$N_i=\exp(-\alpha_T Z_i\psi), \tag{10}$$

$$\partial u_D/\partial\tau+u_D\partial u_D/\partial\xi=-(q_De/Q_{Do}\alpha_{aD})\partial\psi/\partial\xi, \tag{11}$$

$$\partial N_D/\partial\tau+(\partial/\partial\xi)(N_Du_D)=0, \tag{12}$$

$$\nabla_\perp^2\psi+\partial^2\psi/\partial\xi^2=(1/\alpha_{ne})(N_e-Z_iN_i-N_Dq_DQ_{Do}/e), \tag{13}$$

$$\begin{aligned}\partial q_D/\partial\tau+u_D\partial q_D/\partial\xi=A_1[&-N_e(T_e/m_e)^{1/2}\exp(\psi_f-\psi_p)\\&+Z_iN_i(T_i/m_i)^{1/2}\\&\times\{1-Z_i\alpha_T(\psi_f-\psi_p)\}],\end{aligned} \tag{14}$$

where $N_i=n_i/n_{io}$, $N_e=n_e/n_{io}$, $N_D=n_D/n_{io}$, $u_D=v_D/\omega_{pD}\lambda_{\rm De}$, $\tau=t\omega_{pD}$, $\xi=x/\lambda_{\rm De}$, $q_D=Q_D/Q_{Do}$, $\psi_f=e\phi_f/k_BT_e$, etc. $\alpha_{ne}=n_{eo}/n_{io}$, $\alpha_{nD}=n_{eo}/n_{Do}$, $\alpha_T=T_e/T_i$, and $A_1=\pi e a^2(8/\pi)^{1/2}$. In the above equations the Laplace operator ∇^2 has been decomposed into a transverse part $\nabla_\perp^2$ and a longitudinal part $\partial^2/\partial\xi^2$. Here we assume that the ξ axis is the direction of propagation of the wave, and is also the axis of the wave guide. To deduce the dispersion relation we assume that the field variables are perturbed as

$$N_e=N_{eo}+N_e',\ \ \psi=\psi',\ \ N_i=N_{io}+N_i',\ \ u_D=u_D',$$

$$N_D=N_{Do}+N_D',\ \ q_D=q_{Do}+q_D'=1+q_D', \tag{15}$$

and linearize Eqs. (9)–(14) to obtain

$$N_e'=\alpha_{ne}\psi', \tag{16}$$

$$N_i'=-\alpha_TZ_i\psi', \tag{17}$$

$$\partial u_D'/\partial\tau=-q_{Do}(e\alpha_{aD}/Q_{\rm De})\partial\psi'/\partial\xi, \tag{18}$$

$$\partial N_D'/\partial\tau=-N_{Do}\partial u_D'/\partial\xi, \tag{19}$$

$$\nabla_\perp^2\psi'+\partial^2\psi'/\partial\xi^2=(1/\alpha_{ne})[N_e'-Z_iN_i'-Q_{Do}/e(N_{Do}q_D'+q_{Do}N_D')], \tag{20}$$

$$\begin{aligned}\partial q_D'/\partial\tau=(I_{eo}/\alpha_{ne})N_e'+I_{io}N_i'+[&(I_{eo}/\alpha_{ne})N_{eo}\\&-A(m_e/m_i\alpha_T)^{1/2}Z_i^2N_{io}]\psi_f',\end{aligned} \tag{21}$$

where $I_{eo}=-A\alpha_{ne}\exp(\psi_{fo}-\psi_p)$, $I_{io}=AZ_i\{(m_e/m_i)/\alpha_T\}^{1/2}\{1-Z_i\alpha_T(\psi_{fo}-\psi_p)\}$, $A=A_1(T_e/m_e)^{1/2}$, and $\psi_f'=q_D'/C$.

Now we assume the perturbed quantities to vary as follows:

$$\begin{aligned}N_e'&=f(r)\exp[i(k\xi-\omega\tau)],\\N_i'&=g(r)\exp[i(k\xi-\omega\tau)],\\N_D'&=h(r)\exp[i(k\xi-\omega\tau)],\\u_D'&=l(r)\exp[i(k\xi-\omega\tau)],\\q_D'&=m(r)\exp[i(k\xi-\omega\tau)],\\\psi'&=n(r)\exp[i(k\xi-\omega\tau)],\end{aligned} \tag{22}$$

from which Eqs. (16)–(21) lead to

$$f(r)=\alpha_{ne}n(r), \tag{23}$$

$$g(r)=-\alpha_TZ_in(r), \tag{24}$$

$$l(r)=(k/\omega)(q_{Do}e\alpha_{nD}/Q_{Do})n(r), \tag{25}$$

$$h(r)=(kN_{Do}/\omega)l(r), \tag{26}$$

$$d^2n(r)/dr^2+(1/r)dn(r)/dr-k^2n(r)=G, \tag{27}$$

where

$$\begin{aligned}G=(1/\alpha_{ne})[&f(r)-Zig(r)-(Q_{Do}N_{Do}/e)m(r)\\&-(Q_{Do}q_{Do}/e)h(r)],\end{aligned} \tag{28}$$

$$\begin{aligned}(I_{eo}/\alpha_{ne})f(r)+I_{io}g(r)=\{&(1/C)[AZ_i^2N_{io}(m_e/m_i\alpha_T)^{1/2}\\&-(I_{eo}/\alpha_{ne})N_{eo}]-i\omega\}m(r).\end{aligned} \tag{29}$$

Eliminating all the variables in favor of $n(r)$, we obtain

$$\begin{aligned}d^2n(r)/dr^2+(1/r)dn(r)/dr+k^2[&\{(1/\omega^2)(N_{Do}\alpha_{nD}/\alpha_{ne})\\&+(1/k^2)(Q_{Do}N_{Do}/eb\alpha_{ne}-Z_i^2\alpha_T/\alpha_{ne}-1)\}-1]n(r)\\=0.\end{aligned} \tag{30}$$

If we put

$$\begin{aligned}\beta^2=k^2[&\{(1/\omega^2)(N_{Do}\alpha_{nD}/\alpha_{ne})+(1/k^2)(Q_{Do}N_{Do}/eb\alpha_{ne}\\&-Z_i^2\alpha_T/\alpha_{ne}-1\}-1],\end{aligned} \tag{31}$$

then the solution of Eq. (30) is seen to be

$$n(r)=J_o(\beta r). \tag{32}$$

Now, on the boundary of the cylindrical wave guide we should have $\psi=0$, that is,

$$J_o(\beta R)=0, \tag{33}$$

if R is the radius of the cylinder. The dispersion relation then turns out to be

$$p_{on}^2=(KR)^2[\{(1/\omega^2)(N_{Do}\alpha_{nD}/\alpha_{ne})+(1/k^2)\times(Q_{Do}N_{Do}/eb\alpha_{ne}-Z_i^2\alpha_T/\alpha_{ne}-1)\}-1], \tag{34}$$

p_{on} being the root of Eq. (33).
Here $b=[(1/C)\{a^2eZ_i^2T_e(8\pi/m_iT_i)^{1/2}-I_{eo}\}-i\omega]/(I_{eo}-I_{io}Z_i\alpha_T)$, where

$$C=a(1+a/\lambda_D) \tag{35}$$

is the capacitance of the dust grain. We have solved the dispersion relation for various values of the radius R and three different zeroes of the Bessel function. The corresponding values of the phase velocity and imaginary wave vector (k_{im}) are plotted, respectively, in the Figs. 1(a), 1(b), 2(a), and 2(b). In each case the variation of the phase velocity and k is depicted with respect to the frequency ω. It may be observed that for $R=10$ and larger the behavior is almost the same as an infinite plasma.

III. NONLINEAR ANALYSIS: SAGDEEV POTENTIAL

To study the nonlinear regime of the above system of equations we take recourse to the pseudopotential approach. This method, being applicable even for large amplitude waves, will yield a positive result regarding the existence of a nonlinear solitary wave in the system. We again start with a perturbation of the type of Eq. (15), but do not linearize or assume the perturbed quantities to be of the form of Eqs. (16)–(21); we obtain

$$N_e'=\alpha_{ne}\psi', \tag{36}$$

$$N_i'=-\alpha_T Z_i\psi', \tag{37}$$

$$\partial u_D'/\partial t+u_D'\partial u_D'/\partial x=-(1+q_D')(e\alpha_{nD}/Q_{Do})\partial\psi'/\partial x, \tag{38}$$

$$\partial N_D'/\partial t+N_{Do}\partial u_D'/\partial x=-\partial(u_D'N_D')/\partial x, \tag{39}$$

$$\partial^2\psi'/\partial r^2+(1/r)\partial\psi'/\partial r+\partial^2\psi'/\partial x^2=(1/\alpha_{ne})[N_e'-Z_iN_i'-(Q_{Do}/e)(N_D'+q_D'N_D'+N_{Do}q_D')], \tag{40}$$

$$\partial q_D'/\partial t=(I_{eo}/\alpha_{ne})(N_e'+q_D'N_e'/C+N_{eo}q_D'/C)+I_{io}N_i'-A(m_e\alpha_T/m_i)Z_i^2/C(q_D'+N_i'q_D'), \tag{41}$$

Here x and t are the normalized length and time. Actually, there are τ and ξ in place of t and x in the mother equations,

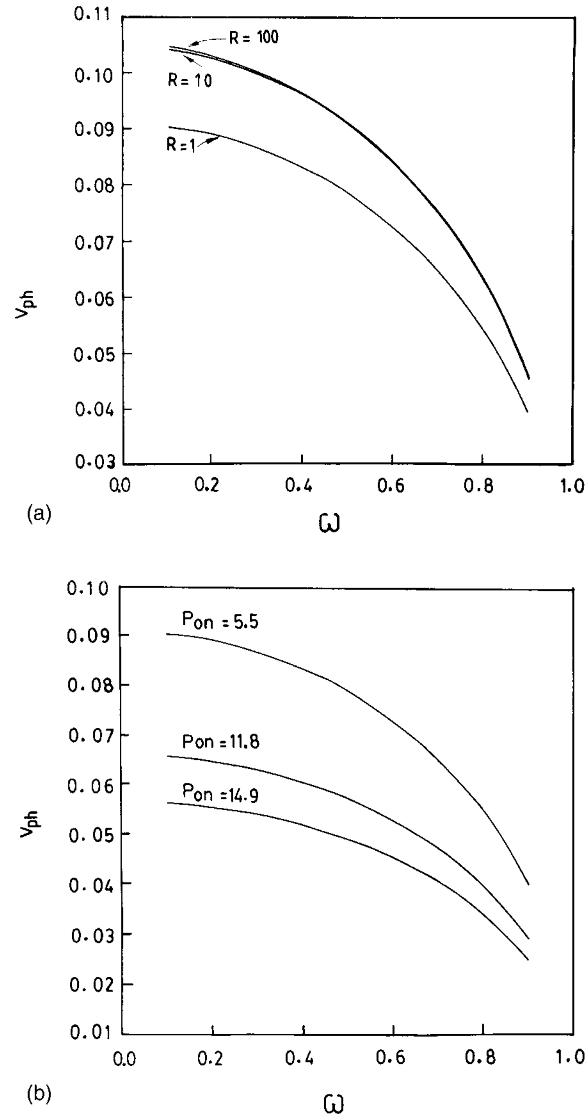

FIG. 1. (a) Variation of V_{ph}(cm/sec), with ω, where V_{ph} is the phase velocity of the dust acoustic wave for various radii of the cylindrical wave guide (R) when $p_{\mathrm{on}}=5.5$. (b) Variation of V_{ph}(cm/sec), with ω, for different values of p_{on}, the root of $J_0(\beta R)=0$ when $R=1$ cm.

and we have omitted the convective term for the dust charge fluctuation. We now assume that the perturbed quantities (u_D',ψ',N_i',N_D', etc.) all have the general variations:

$$\begin{aligned} N_e'&=J_o(k_\perp r)N_e(x,t),\\ N_i'&=J_o(k_\perp r)N_i(x,t),\\ N_D'&=j_o(k_\perp r)N_D(x,t),\\ u_D'&=J_o(k_\perp r)u_D(x,t),\\ q_D'&=J_o(k_\perp r)q_D(x,t),\\ \psi'&=J_o(k_\perp r)\psi(x,t), \end{aligned} \tag{42}$$

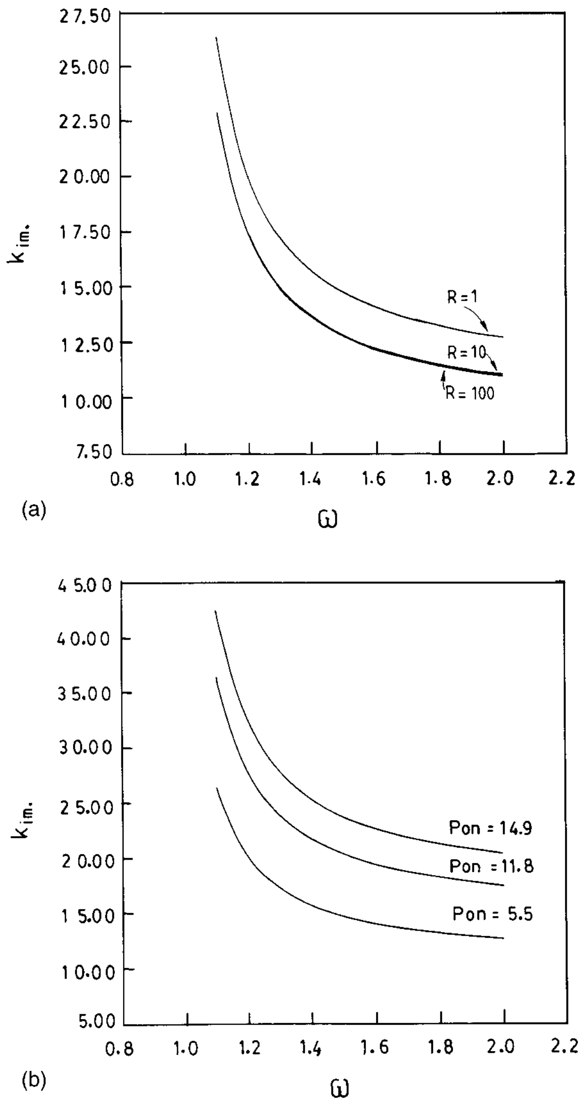

FIG. 2. (a) Variation of $k_{\rm im}$, the imaginary propagation constant, as obtained from the dispersion relation with ω for various radii (R) when $p_{\rm on}=5.5$. (b) Variation of $k_{\rm im}$ with ω for different values of $p_{\rm on}$ when $R=1$ cm.

Substituting Eq. (42) into Eqs. (36)–(41), we obtain

$$N_e=\alpha_{ne}\psi, \tag{43}$$

$$N_i=-\alpha_T Z_i\psi, \tag{44}$$

$$\partial u_D/\partial t+\alpha u_D\partial u_D/\partial x=-(e\alpha_{nD}/Q_{Do})\partial\psi/\partial x -q_D(e\alpha_{nD}/Q_{Do})\alpha\partial\psi/\partial x, \tag{45}$$

$$\partial N_D/\partial t+N_{Do}\partial u_D/\partial x+\alpha\partial/\partial x(u_D N_D)=0, \tag{46}$$

$$\partial q_D/\partial t=(I_{eo}/\alpha_{ne})(N_e+(1/C)\alpha q_D N_e+N_{eo}q_D/C)+I_{io}N_i -AZ_i^2(m_e\alpha_T/m_i)^{1/2}(1/C)(\alpha N_i q_D+q_D), \tag{47}$$

where

$$\alpha=\int_O^R rJ_o^3 dr \Big/ \int_O^R rJ_o^2 dr, \tag{48}$$

Now, following Ma and Liu, we note that the time scales for grain charging and dust hydrodynamic motion are widely different. The charging time is approximately

$$\tau_{ch}=[(dQ_D/dt)(1/Q_D)]^{-1}\sim Q_{Do}/I_o, \tag{49}$$

while the hydrodynamic time scale is of the order of $\tau_d\sim\omega_{pd}^{-1}=(4\pi Q_{Do}^2 N_{Do}/m_D)^{-1/2}$, where N_{Do} and $Q_{Do}\sim-2.5aT_e/e^7$ are dust density and charge in equilibrium, and

$$I_o=-\pi a^2 e(8T_e/\pi m_e)^{1/2}n_{eo} \tag{50}$$

is the electron grain current evaluated at $q_D=0$. For a dusty plasma of $n_{eo}\sim10^{10}$ cm^{-3} and $T_e\sim1$ eV as well as 1-μm grains of $N_{Do}\sim10^5$ cm^{-3}, $m_D\sim4\times10^{-12}$ g, $\tau_{\rm ch}=95$ ns, $\tau_d=2$ ms So $\tau_{\rm ch}\ll\tau_d$ that is on the hydrodynamic time scale the dust charge can reach equilibrium at which

$$I_e+I_i=0. \tag{51}$$

So, going to a frame of reference given by $\xi=x-Mt$, we obtain

$$-M\partial u_D/\partial\xi+\alpha u_D\partial u_D/\partial\xi=-(e\alpha_{nD}/Q_{Do})\partial\psi/\partial\xi -\alpha(e\alpha_{nD}/Q_{Do})q_D\partial\psi/\partial\xi, \tag{52}$$

$$N_D=N_{Do}u_D/(M-\alpha u_D), \tag{53}$$

$$\partial^2\psi/\partial\xi^2-\psi=(1/\alpha_{ne})[N_e-Z_iN_i-(Q_{Do}/e)(N_D+\alpha q_D N_D +N_{Do}q_D)], \tag{54}$$

$$N_e=\alpha_{ne}\exp(\psi), \tag{55}$$

$$N_i=\exp(-\alpha_T Z_i\psi). \tag{56}$$

On the other hand, $I_e+I_i=0$ leads to

$$q_D+[a_1-a_2\exp(b\psi)]/[a_3+a_4\exp(b\psi)], \tag{57}$$

where

$$\begin{aligned} a_1&=k_B Z_i a/eQ_{Do},\\ a_2&=k_B aB/eQ_{Do},\\ a_3&=Z_i^2/T_i,\\ a_4&=B/T_e,\\ \delta_1&=e\alpha_{nD}/Q_{Do},\\ \delta_2&=\alpha\delta_1,\\ b&=(1+\alpha_1 Z_i),\\ B&=\alpha_{ne}(\alpha_T m_i/m_e)^{1/2}. \end{aligned} \tag{58}$$

Substituting these forms of q_D into Eq. (26) we obtain

$$\alpha u_D^2 - 2Mu_D + \sigma = 0, \tag{59}$$

$$\sigma = 2\delta_1\psi + 2\delta_2 a_2/a_3 \ln\{a_4 \exp(b\psi)\} + 2\delta_2(a_1/a_3 + a_2/a_4)\ln\{a_3 + a_4 \exp(b\psi)\}. \tag{60}$$

Thus we arrive at

$$u_D = (1/\alpha)[M \pm (M^2 - \alpha\sigma)^{1/2}]. \tag{61}$$

Substituting all these into Poisson's equation, we obtain

$$\begin{aligned} d^2\psi/d\xi^2 = {} & \psi + 1/\alpha_{ne}[\alpha_{ne}\exp(b\psi) - Z_i \exp(-\alpha_T Z_i \psi) \\ & - (Q_{Do}/e)\{N_{Do}u_D/(M - \alpha u_D) \\ & + \alpha[N_{Do}u_D/(M - \alpha u_D)] \\ & \times[\{a_1 - a_2\exp(b\psi)\}/\{a_3 + a_4\exp(b\psi)\}] \\ & + N_{Do}[a_1 - a_2\exp(b\psi)]/[a_3 + a_4\exp(b\psi)]\}], \end{aligned} \tag{62}$$

where u_D is given by Eq. (61). Integrating Eq. (62) and using the condition $\partial\psi/\partial\xi \to 0$ as $|\xi| \to \infty$, we obtain the pseudopotential equation

$$\tfrac{1}{2}(\partial\psi/\partial\xi)^2 + V(\psi) = 0, \tag{63}$$

where $V(\psi)$ is given by

$$\begin{aligned} V(\psi) = {} & [\{(Q_{Do}\alpha A_2 N_{Do})/(e\alpha_{ne}M^2 b a_4^2)\}(a_1a_4 + a_2a_3) - (1 - Z_i/\alpha_{ne} - Q_{Do}N_{Do}/eM^2\alpha_{ne}) - (N_{Do}Q_{Do}/e\alpha_{ne})(a_2/a_4) \\ & \times(1 + \alpha A_1/M^2)]\psi - [1 + Z_i^2\alpha_T/2\alpha_{ne} - (A_2Q_{Do}N_{Do})/(2e\alpha_{ne}M^2) + (Q_{Do}N_{Do}\alpha A_2/2eM^2\alpha_{ne})(a_2/a_4)]\psi^2 \\ & + [\{(N_{Do}Q_{Do})/(eba_4^2\alpha_{ne})\}(a_1a_4 + a_2a_3)(1 + \alpha A_1/M^2 - \alpha A_2 a_5/M^2 b a_4)]\ln(a_5 + ba_4\psi) \\ & - [N_{Do}Q_{Do}(a_1a_4 + a_2a_3)/(e\alpha_{ne}ba_4^2)](1 + \alpha A_1/M^2 - \alpha A_2 a_5/M^2 b a_4)\ln a_5, \end{aligned} \tag{64}$$

where

$$A_1 = \delta_2[(a_1 \ln a_4)/a_3 + (a_1/a_3 + a_2/a_4)\ln(a_3 + a_4)],$$

$$A_2 = [\delta_1 + b\delta_2\{a_1/a_3 + (a_1/a_3 + a_2/a_4)a_4/(a_3 + a_4)\}],$$

$$a_5 = a_3 + a_4.$$

a_1, a_2, a_3, a_4, b, and B were defined earlier.

IV. ANALYSIS AND DISCUSSION

From the form of Eq. (63), it is quite obvious that it is impossible to obtain the form of ψ analytically. So we analyze the form of potential $V(\psi)$ numerically for different plasma parameters. The corresponding forms of $V(\psi)$ are depicted in the Figs. 3(a)–3(c). The parameter values taken are $\alpha_T = 100$ and $\alpha = -4.9$, -25.35, and -40.88, whereas $\alpha_{nD} = 100$, 200, and 300. The Mach numbers are 0.5, 1, and 1.5. The form of the pseudopotential obtained shows that the particle trapping is not possible. So a solitary wave can not be formed. Recent theoretical investigations by various authors (see Refs. [15], [28–30]) revealed that a plasma contaminated with dust grains of constant charge may and may not sustain a soliton (dust acoustic) depending on the techniques followed Ref. [31], the models used Refs. [31,32], the range of values of the Mach number, etc. It was reported by some workers [10] that if dust dynamics and a self-consistent dust grain charge fluctuation be incorporated, then the propagation of a large-amplitude dust acoustic soliton through a duty plasma is possible, *within a certain range of velocities or Mach numbers*. This implies that a soliton will not exist in the plasma if the Mach number is beyond this range. In our paper, the nontrapping solution does not favor soliton formation. This may be due to either the values of the Mach numbers used in the numerical computation or due to a nonlinear dissipative mechanism introduced by the nonsteady dust charge fluctuation dynamics. This inference is also supported by the analysis of Ref. [33], where by using reductive perturbation theory, the authors deduced a Burger equation. It is a well known fact that the Burger equation does not support a soliton solution. The difference between the treatment of Ref. [33] and our treatment is that for a bounded system, their analysis is valid for an infinite plasma. At this point it may be pointed that a reductive perturbation approach in a bounded domain is more difficult than the treatment given here.

V. CONCLUSION

In our analysis we have theoretically investigated the propagation of large amplitude dust acoustic waves in a bounded dusty plasma system, as well as the existence of solitary waves in the plasma. It is seen that the presence of charged dust particles in a plasma gives some fascinating characteristics of the potential profile in comparison with those found in other kinds of plasma. In the present analysis we find that a major portion of the Sagdeev potential is positive in the presence of dust charges in the plasma. The peculiar behavior of the Sagdeev potential indicates that the possibility of the existence of a dust acoustic soliton in a bounded dusty plasma is almost nil; instead, double layers

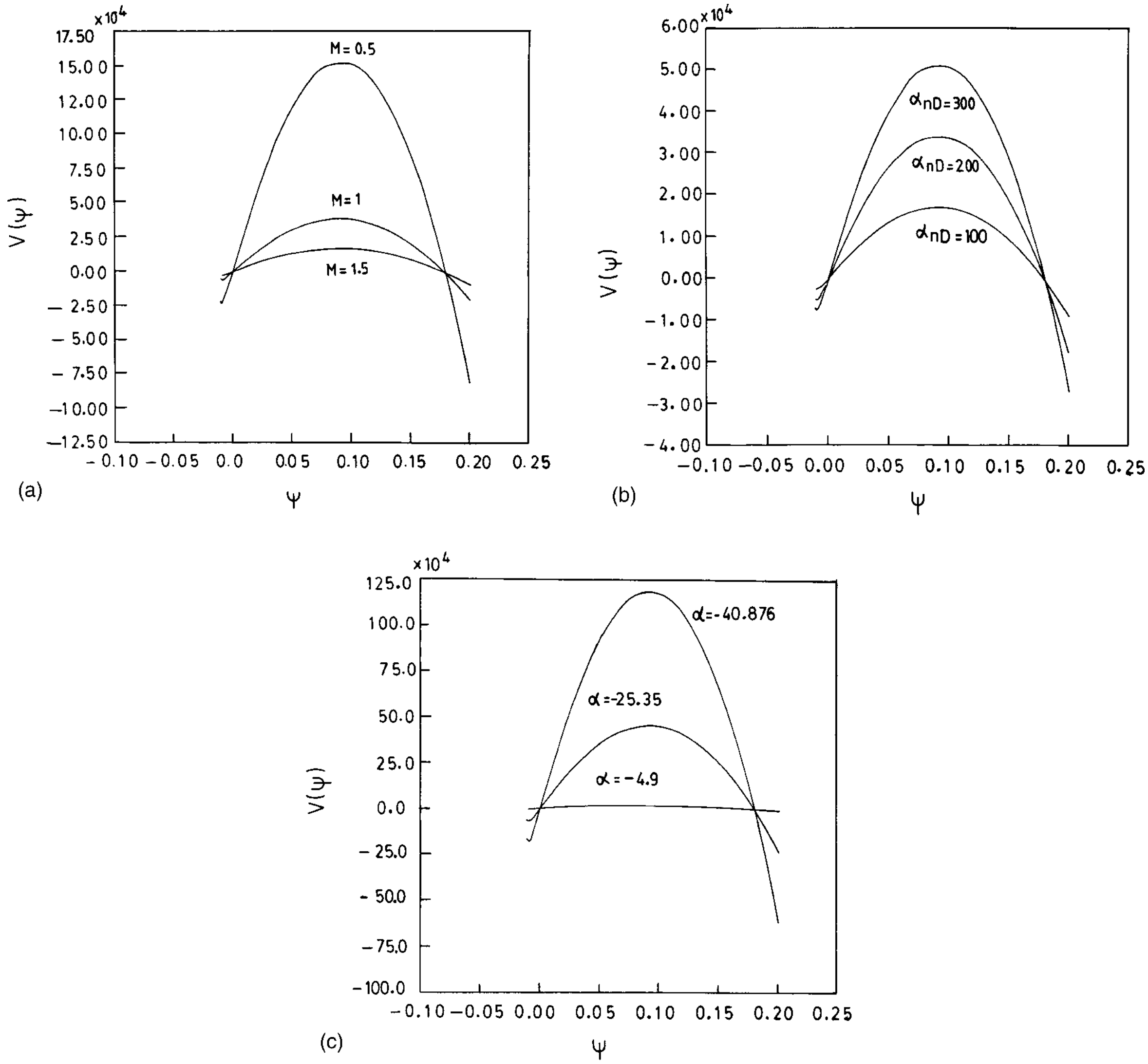

FIG. 3. (a) Plot of $V(\psi)$ against ψ, where V is the Sagdeev potential for various values of M, and $\alpha=-4.9$, corresponding to $p_{\mathrm{on}}=5.5$ when the typical plasma parameters are $\alpha_{nD}=n_{eo}/n_{Do}=100$, $\alpha_T=T_e/T_i=100$ ($T_e=100$ eV, $T_i=1$ eV), and $Q_{Do}=10e$. (b) Variation of $V(\psi)$ with ψ for different values of α_{nD} when $\alpha=-4.9$ and $M=1.5$. (c) Variation of $V(\psi)$ with ψ for various values of α corresponding to different values of p_{on} when $\alpha_{nD}=100$ and $M=1.5$.

will occur. However, we have not compared our theoretical results with experimental observations due to the nonavailability of experimental data, although much work was done in Refs. [34–38] and others in studying the existence of solitary waves in laboratory plasmas. In the present study of the existence of solitary waves in a bounded dusty plasma, the dust fluid has been assumed to be cold. The effect of dust on temperature may be studied in order to make our analysis more realistic from an experimental point of view. Moreover, dust grains have been considered to be spheres of equal radius that carry identical charges. But this is not the case for a real dusty plasma, which consists of dust grains with size and charge distributions. It would be better to incorporate these features for a real situation.

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