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Adiabatic Electron Acceleration in a Cnoidal Wave

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Abstract

Acceleration and energy gain of a relativistic charged particle in the field of a cnoidal wave with slow variation in amplitude is discussed. It is shown that solitary waves are more effective in imparting energy to a charged particle than the usual sinusoidal waves.

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An efficient way of particle acceleration is to exploit the large electric fields associated with longitudinal electrostatic waves. Since the phase velocity and the accelerating electric field of these waves are in the same direction, it is possible to 'phase-lock' a beam of injected particles and accelerate them to very high energies. The large amplitude electric fields neccesary for such a purpose can arise naturally in plasmas owing to its ability of supporting space charge waves with relativistic phase velocities. This concept has lead to extensive theoretical investigations on plasma-based accelerators such as plasma beat wave [1,2] and wake field [3,4] accelerators.

While studying the problem of acceleration of particles by waves, a realistic approach would be to consider a localized wave packet where the amplitude of the wave has a slow spatial variation. If the flight time of a particle across the wave packet is much greater than the bouncing time, the particle will be trapped within the potential field of the wave. In the phase space, trajectories of trapped particles correspond to closed curves lying within a region bounded by separatrices, while those of untrapped particles correspond to open curves lying outside the separatrix. For a localized wave packet, particles moving in the wave field experience a slow variation in wave amplitude. When the amplitude of the wave increases, the trajectory of an untrapped particle in phase space crosses the separatrix to enter the trapped region and when the amplitude of the wave decreases, the phase space trajectory once again crosses the separatrix so that the particle becomes detrapped. To analyze the motion of particles under such a situation, one should conveniently take a recourse to the adiabatic theory and formulate such wave-particle interaction processes. When the particle orbit in phase space crosses a separatrix the adiabatic assumption is violated because of the divergence of particle's bounce period. The change in the adiabatic invariant associated with the crossing of the separatrices by the particle trajectory in phase space can be related to the corresponding change (gain or loss) of the particle energy.

The acceleration of charged particles both in the nonrelativistic and relativistic sinusoidal waves in the adiabatic limit has been studied both analytically and numerically [5 - 7]. However, such investigations in the fields of waves other than sinusoidal are rather scanty. Since nonlinearity in wave propagation is an important occurence in physical processes of the present day plasma physics, particularly in the context of large amplitude waves being considered [8], it is rather imperative to investigate explicitly the energy exchange between the charged particle and a wave which is governed by a nonlinear equation. In this paper we have developed the adiabatic theory of electron acceleration in a relativistic cnoidal wave and a solitary wave with sech^2 type profile.

The wave field solutions of several nonlinear equations such as nonlinear Schrodinger equation can be modelled [9] by the following expression

$$y = A\ cn(\xi \mid \kappa) \tag{1}$$

The symbol $cn(\xi \mid \kappa)$ stands for a Jacobi elliptic function [10] with $\xi = x - vt$, and κ is a parameter with $0 \leq \kappa \leq 1$. The periodicity length of y as a function of variable ξ is given by $4K(\kappa)$, where K is a complete elliptic integral of first kind. Two important limits of the above solution are

$$y \rightarrow A\ \cos\xi \ \text{ and } \ K(\kappa) \rightarrow \pi/2 \ \text{ for } \ \kappa \rightarrow 0$$

which corresponds to the usual sinusoidal wave and

$$y \rightarrow \frac{A}{\cosh\xi} \ \text{ and } \ K(\kappa) \rightarrow \infty \ \text{ for } \ \kappa \rightarrow 1$$

which is a soliton like solution.

The relativistic Hamiltonian of an electron interacting with the wave is given by

$$H(q,p,t) = \sqrt{p^2c^2 + m_e^2c^4} - m_ec^2 + A(q)\ cn(kq - \omega t \mid \kappa) \tag{2}$$

where, m_e is the mass of the electron, p, q are its momentum and position respectively. $A(q)$ is the slow space dependent amplitude of the electrostatic potential, $A = keE/m_e\omega^2$ where e is the charge of the electron, k is the wave number and ω is the wave frequency.

To investigate the motion of the electron in a wave packet with slowly varying amplitude in the adiabatic theory [11], we make transformations, as in Bruhwiler and Cary [6], so that q becomes a 'temporal coordinate' and $\theta = kq - \omega t$ (position in the wave frame), becomes a position coordinate. With this transformation, the new momentum p_θ and the corresponding Hamiltonian H_θ become,

$$p_\theta = \frac{H}{\omega}, \quad H_\theta = \frac{k}{\omega}H - p = kp_\theta - p \tag{3}$$

From eqs. (2) and (3) we get,

$$p = \pm\sqrt{\frac{1}{c^2}[\omega p_\theta - A\ cn(\theta \mid \kappa) + m_e c^2]^2 - m_e^2 c^2} \tag{4}$$

so that,

$$H_\theta = kp_\theta \mp \sqrt{\frac{1}{c^2}[\omega p_\theta - A\ cn(\theta \mid \kappa) + m_e c^2]^2 - m_e^2 c^2} \tag{5}$$

In the following, we consider the minus sign in Eq. (5) which corresponds to particles going in the same direction as the wave.

To determine the separatrices in the $\theta - p_\theta$ phase space, we first obtain the fixed points from the following equations:

$$\frac{d\theta}{dq} = \frac{\partial H_\theta}{\partial p_\theta} = 0, \quad \frac{dp_\theta}{dq} = -\frac{\partial H_\theta}{\partial \theta} = 0. \tag{6}$$

So the saddle points in phase space, as obtained from eqn(6), are given by:

$$sn(\theta \mid \kappa) = 0, \text{i.e. at } \ \theta = 0, 4K, \ldots \tag{7}$$

The values of momenta $p_{\theta s}$ and Hamiltonian $H_{\theta s}$ at the saddle points defined by eq. (7) are given by the following expressions:

$$p_{\theta s} = \frac{m_e\omega}{k^2}\frac{\Gamma_p^2}{\Gamma_p+1} + \frac{A}{\omega} \tag{8a}$$

$$H_{\theta s} = -\frac{m_e\omega}{k}\frac{\Gamma_p}{\Gamma_p+1} + \frac{kA}{\omega} \tag{8b}$$

where, $\Gamma_p^2 = 1/(1-\beta_p^2)$ and $\beta_p = \omega/kc$, is the relativistic factor associated with the wave phase velocity ω/k.

From the above expressions it is explicit that the saddle points (and the separatrices) are functions of the slow 'time' variable q. In order to make way for the application of the separatrix crossing theory [12] we should make a canonical transformation which would render the fixed points 'stationary' (i.e. independent of 'time' q) in phase space. Following Mora [7], a canonical transformation from (θ, p_θ) to (ϕ, p_ϕ) such that $p_\phi = 0$ at the saddle points and $H_\phi = 0$ on the separatrices is obtained through the generating function of F_2 type:

$$\begin{aligned} F_2(p_\phi, \theta, q) &= p_\phi\theta + \frac{m_e\omega}{k^2}\frac{\Gamma_p^2}{\Gamma_p+1}\theta \\ &+ \frac{A}{\omega}\int^\theta cn(\theta \mid \kappa)d\theta + \int^q \left(\frac{m_e\omega}{k}\frac{\Gamma_p}{\Gamma_p+1} - \frac{kA}{\omega}\right) dq' \end{aligned} \tag{9}$$

The Hamiltonian H_ϕ in the new (ϕ, p_ϕ) coordinates may be written after some algebra as

$$H_\phi = kp_\phi - \frac{k}{\omega}B + \frac{m_e\omega\Gamma_p}{k} - \sqrt{\frac{m_e^2\omega^2}{k^2}\Gamma_p^2 + 2m_e\omega\Gamma_p p_\phi + k^2\beta_p^2 p_\phi^2} \tag{10}$$

with $B = A[1 - cn(\phi \mid \kappa)]$. We can then obtain the expressions for $(p_\phi)_\pm$, the values of p_ϕ at the upper and lower branches of the separatrix by putting $H_\phi = 0$ and solving for p_ϕ. We get

$$(\omega p_\phi)_\pm = \Gamma_p^2 \left[B \pm \sqrt{\beta_p^2 B^2 + 2\frac{m_e\omega^2}{k^2}\frac{B}{\Gamma_p}} \right] \tag{11}$$

If the amplitude of the wave is a constant, the motion of an electron in the $(\phi - p_\phi)$ phase space is confined to lines of constant Hamiltonian. Due to the dependence of the amplitude on time, the Hamiltonian is no longer an integral of motion and hence the phase trajectories also change in time. However, due to the slow variation of amplitude with time, transition from one trajectory to another is governed by the conservation of the action integral for the untrapped particles. The action integral is defined by

$$J_\pm = \int_0^{4K} (p_\phi)_\pm \, d\phi \tag{12}$$

where $(p_\phi)_\pm$ is defined by eq. (11). It is convenient to express the electron energy in terms of the relativistic expression Γ defined by $\Gamma = (\sqrt{p^2c^2 + m_e^2c^4})/m_ec^2$. so that.

$$\Gamma_\pm = \Gamma_p + \frac{\Gamma_p^2}{4K}\int_0^{4K} \left[B \pm \beta_p\sqrt{B^2 + 2\frac{B}{\Gamma_p}} \right] d\phi \tag{13}$$

where the wave amplitude is appropriately normalized. For the case of solitary waves, we take the wave electrostatic potential in eq. (2) as $A(q)\ \mathrm{sech}^2(kq - \omega t)$, so that the upper and lower bounds of electron trapping energy given in eq. (13) can be worked out exactly to yield

$$\Gamma_\pm = \Gamma_p + \Gamma_p^2 A \pm \Gamma_p^2 \beta_p \sqrt{A^2 + \frac{2A}{\Gamma_p}} \tag{14}$$

So for the electron to be trapped, the upper and lower bounds of Γ for a given value of the maximum amplitude of the wave A_0 are

$$\Gamma_-(A_0) < \Gamma < \Gamma_+(A_0) \tag{15}$$

where $\Gamma_\pm(A_0)$ is the value of $\Gamma_\pm$ at the maximum amplitude of the wave. Those particles whose initial energies Γ_i are such that $\Gamma_i < \Gamma_-(A_0)$ or $\Gamma_i > \Gamma_+(A_0)$ do not interact and exchange energy with the wave.

We would like to calculate the energy gain (loss) of an electron which enters the region bounded by the separatrices from below (above) and finally leaves that region from above (below). In these conditions, the energy variation mechanism is related to the breaking of the adiabatic invariant as the particle trajectory crosses the separatrix. The amplitude $A = A_x$, at the instant at which the particle is trapped, is such that

$$J_\pm(A_x) = \frac{4K}{\omega}(\Gamma_i - \Gamma_p)$$

When the amplitude increases from A_x to its maximum value A_0, the particle remains trapped. The particle is again detrapped when the amplitude decreases from A_0 to 0 at the same value of $A = A_x$. It is evident from eqs. (13) and (14) that the energy of the particle, Γ, consists of two branches. A net energy variation of the electron can occur if during the course of transition, the trajectory starting from branch $\Gamma = \Gamma_i$ ends on a different branch $\Gamma = \Gamma_f$, where Γ_f and Γ_i are two solutions for Γ. This implies a double crossing of the separatrix at the same value of $A = A_x$. The adiabatic invariant breaks at each crossing of the separatrix and remains constant during each stage of motion. The energy variation of the particle during a transition is then given by the variation of the action integral as $\Delta\Gamma = \omega\Delta J/4K$. Using eqs. (11) and (12) we get

$$\Delta\Gamma = \frac{\omega(J_+ - J_-)}{4K} = \frac{2\Gamma_p^2\beta_p}{4K}\int_0^{4K}\sqrt{B^2 + 2\frac{B}{\Gamma_p}}\, d\phi \qquad (16a)$$

and for the solitary wave

$$\Delta\Gamma = 2\Gamma_p^2\beta_p\sqrt{A^2 + 2A/\Gamma_p} \qquad (16b)$$

The energy change averaged over initial electron phases can be expressed as $\langle\Delta\Gamma\rangle = P\Delta\Gamma$, where P represents the probability of transition to different regions of phase space. The probability [13] can be computed in terms of time variation of the area of phase space regions bounded by the separatrix determined at the instant when the crossing of the separatrix occurs for a uniform distribution of particles in the phases of motion. Eqs. (16) give the energy change of an electron of a given energy during a transition of the wave packet.

We next proceed to determine the average energy gain of an electron. The average absorbed energy normalized by $m_e c^2$ can be written in the following form

$$E = \int_{\Gamma_-(A_0)}^{\Gamma_+(A_0)} P \Delta\Gamma f(\Gamma) \frac{\Gamma d\Gamma}{\sqrt{\Gamma^2 - 1}} \tag{17}$$

where $f(\Gamma)$ is the one dimensional relativistic distribution function (normalized to unity)

$$f(\Gamma) = \frac{1}{2\mathcal{K}_2(\mu)\mu} (1 + \mu\, \Gamma)\, e^{-\mu\Gamma}$$

where $\mu = m_e c^2 / k_B T_e$ and $\mathcal{K}_2(\mu)$ denotes a modified Bessel function of the second kind with index 2 and argument μ. We note that the trapping condition as well as the absorbed energy E depend only on the maximum amplitude of the wave electric field and not on its profile, provided that the adiabaticity condition for the particle motion is met.

For a given value of the maximum field amplitude A_0, the energy variation $\Delta\Gamma$ is non zero only for a set of initial energies defined by eq. (15). By means of eq. (17) the average energy gain of an electron is reduced to an integration. The equation gives the average absorbed energy by means of three essential ingredients : the range of energy values arising out of the trapping condition given in eq. (15), the energy change during a transition which depends on the initial energy of the particle (energy change $\Delta\Gamma$ depends implicitly on the initial energy as dictated by eqs. (13),(14) and (16)), and the transition probability P. The integral given in eq. (17) is evaluated numerically. Fig. 1 shows some results for the average energy gain E for $\mu = 2.0$ and normalized value of $A_0 = 0.1$ as a function of Γ_p for different values of κ and for the sech^2 type wave packet.

It is found that for a given Γ_p, the energy gain increases with κ and thus the cnoidal wave is more efficient in transferring energy to a charged particle. However, for a solitary wave, the gain of energy by a particle is much higher than that in a cnoidal wave (even when $\kappa = 0.9$). The singular behaviour of the solitary wave in transferring energy to a charged particle can be explained in the following way:

As stated above, the average energy gained by a charged particle depends on three factors ; i) the trapping condition giving the range of initial energy for which the particle will be trapped, ii) the energy change during a transition, ($\Delta\Gamma$) and iii) the transition probability. For a solitary wave (in the limit $\kappa \to 1$), the trapping condition defined by eqn (14) and (15) produces an increase in the range of trapping energies and thus causes a larger number of particles to participate in the wave particle interaction (through the factor $e^{-\mu\Gamma}$). A comparison between the range of energies for particles participating in the wave particle interaction can be made from eqn(13) and (14). For cnoidal wave ($\kappa = 0.9$), $\Gamma_- = 2.4$ and 5.1 for $\Gamma_p = 5$ and 20 respectively; whereas for solitary wave the corresponding values are 2.0 and 3.5. The lower values of Γ_- causes the numerical value of the probabilty distribution factor $e^{-\mu\Gamma}$ to increase, for solitary wave; (a two fold increase at $\Gamma_p = 5$ and an almost tenfold increase at $\Gamma_p = 20$). Physically, this indicates that more number of particles (with lower initial energies) are involved in the wave particle interaction for solitary wave.

Moreover, for solitary waves with larger values of Γ_p, the decrease in the probability distribution function $e^{-\mu\Gamma}$ (due to an increase of Γ_-, the lower bound of trapping energy) is adequately compensated by the increase in the single particle energy gain $\Delta\Gamma$, so that the average energy gain does not vary much with Γ_p. On the other hand, for cnoidal waves, the number of particles interacting with the wave with energies in the neighbourhood of Γ_- is too small for larger value of Γ_p, (as Γ_- attains a higher value for higher Γ_p) so the average energy gain decreases with the increase of Γ_p.

However, an important point regarding the validity of adiabatic approximation for the solitary wave should be made explicit. As $\kappa \to 1$, the periodicity length increases causing the bounce frequency ω_b of the trapped particle to decrease and the adiabatic approximation breaks down for a wave packet of solitary wave. In this case the merit of this investigation could perhaps be advocated by stating that this anlysis of the energy exhange between a charged particle and a solitary wave (where the solitary wave can be regarded as a limiting case of a cnoidal type of wave) gives a basic understanding regarding the efficiency of different spatially adibatic waves in transferring energy to particles. Moreover, in a realistic situation, collisions, fluctuations or noise present in the system scatter resonant particles more rapidly than ω_b^{-1}, so that the trapped particles are scattered out before a complete period of oscillation. So these results are relevant for the interaction of charged particles with spatially varying large amplitude wave packets.

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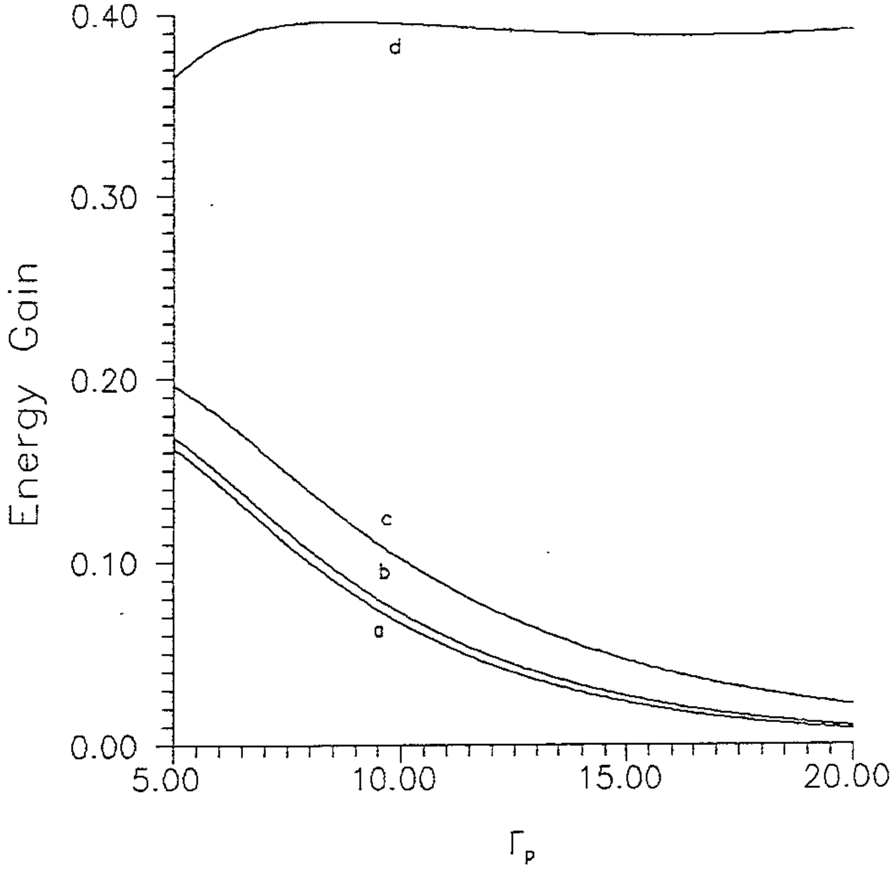
Figure Captions

Fig. 1. Average energy gain of an electron versus Γ_p for (a) κ = 0.0 (b)κ = 0.5 (c) κ = 0.9 (d) sech^2 wave with normalized amplitude of the wave $A_0 = 0.1$.



0.40
0.30
0.20
0.10
0.00
Energy Gain
d
c
b
a
5.00
10.00
15.00
20.00
Γ_p

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