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Naval Research Laboratory

Washington, DC 20375-5000

NRL Memorandum Report 6224

Stochastic Electron Detrapping in FELs Caused by Sidebands

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July 14, 1988

SECURITY CLASSIFICATION OF THIS PAGE

REPORT DOCUMENTATION PAGE

Form Approved OMB No. 0704-0188

1a REPORT SECURITY CLASSIFICATION UNCLASSIFIED	1b RESTRICTIVE MARKINGS
2a SECURITY CLASSIFICATION AUTHORITY 2b DECLASSIFICATION / DOWNGRADING SCHEDULE	3 DISTRIBUTION / AVAILABILITY OF REPORT Approved for public release; distribution unlimited.
4 PERFORMING ORGANIZATION REPORT NUMBER(S) NRL Memorandum Report 6224	5 MONITORING ORGANIZATION REPORT NUMBER(S)
6a NAME OF PERFORMING ORGANIZATION Naval Research Laboratory 6b OFFICE SYMBOL (If applicable) Code 4790	7a NAME OF MONITORING ORGANIZATION U.S. Army Strategic Defense Command
6c ADDRESS (City, State, and ZIP Code) Washington, DC 20375-5000	7b ADDRESS (City, State, and ZIP Code) Arlington, VA 22245-0280
8a NAME OF FUNDING / SPONSORING ORGANIZATION Strategic Defense Initiative Org. 8b OFFICE SYMBOL (If applicable)	9 PROCUREMENT INSTRUMENT IDENTIFICATION NUMBER
8c ADDRESS (City, State, and ZIP Code) Washington, DC 20301-7100	10 SOURCE OF FUNDING NUMBERS PROGRAM ELEMENT NO 63221C PROJECT NO W31RPD-7-D4039 TASK NO WORK UNIT ACCESSION NO

11 TITLE (Include Security Classification)
Stochastic Electron Detrapping in FELs Caused by Sidebands

12 PERSONAL AUTHOR(S)
Riyopoulos,* S. and Tang, C.M.

13a TYPE OF REPORT	13b TIME COVERED	14 DATE OF REPORT (Year, Month, Day)	15 PAGE COUNT
Interim	FROM ____ TO ____	1988 July 14	51

16 SUPPLEMENTARY NOTATION
*Science Applications Intl. Corporation, McLean, VA 22102

17 COSATI CODES			18 SUBJECT TERMS (Continue on reverse if necessary and identify by block number)
FIELD	GROUP	SUB-GROUP	FELs; Electron detrapping; Sidebands; Efficiency; Stochasticity

19 ABSTRACT (Continue on reverse if necessary and identify by block number)

The growth of sidebands in an FEL above a certain threshold will result in stochastic electron motion. This may lead to significant electron detrapping and loss of amplification for the radiation field. The threshold in the sideband amplitude for the stochastic transition is computed. The rate trapped electrons leak outside the sepratrix is measured by the diffusion coefficient in action space. This approach is general and covers both tapered and untapered wigglers. Three general types of spectra are examined: a narrow, a broad discrete and a broad continuous. Each type is associated with a particular scaling of the diffusion coefficient on the FEL parameters. Numerical results, obtained for constant radiation amplitudes, show good agreement with the theoretical predictions. The diffusion rate is always proportional to the total sideband-to-carrier power ratio, with different coefficients of proportionality for each spectral type. The narrow types of sepctrum causes the

(Continues)

20 DISTRIBUTION / AVAILABILITY OF ABSTRACT ☒ UNCLASSIFIED/UNLIMITED ☐ SAME AS RPT ☐ DTIC USERS	21 ABSTRACT SECURITY CLASSIFICATION UNCLASSIFIED
22a NAME OF RESPONSIBLE INDIVIDUAL C.M. Tang	22b TELEPHONE (Include Area Code) (202) 767-4248 22c OFFICE SYMBOL Code 4790

DD Form 1473, JUN 86 — Previous editions are obsolete. — SECURITY CLASSIFICATION OF THIS PAGE

S/N 0102-LF-014-6603

19. ABSTRACTS (Continued)

highest diffusion rates and the broad continuous spectrum causes the lowest diffusion rates under constant total sideband power. The diffusion length, measured in wiggler periods, is independent of the beam energy γ. Serious deterioration in the FEL efficiency will result when the diffusion length becomes shorter than the wiggler length.

CONTENTS

STOCHASTIC ELECTRON DETRAPPING IN FELS CAUSED BY SIDEBANDS

I. INTRODUCTION

Multifrequency effects in free electron lasers (FEL) become increasingly important as progress is made towards high power operation. Growth of parasitic frequencies (sidebands[1-5]) has been predicted theoretically and has been observed in experiments[6,7] as well as in simulations[8-10] with either constant or tapered wigglers[10]. The efficiency for the carrier signal is reduced and the optical quality is degraded as power is channeled into frequencies apart from the intented operation frequency. Another potential hazard that has attracted little attention so far is the onset of chaotic electron motion caused even by the presence of a single frequency sideband. This may lead to extensive particle detrapping and premature loss of the amplification for all the radiation modes independent of frequency.

So far the attention has been directed to the linear stability issue. The gain for small sideband signal has been computed analytically[1-5] invoking either ensemble averaging over single particle trajectories or solutions of the perturbed kinetic equation for the distribution function. Initial results, obtained for particles localized near the bottom of the ponderomotive well, and, in particular, more recent results including all trapped and untrapped particles[2] with arbitrary distributions, have made it clear that every nontrivial distribution $df_0/dJ \neq 0$ is unstable to sideband growth.

Given that sidebands cannot be eliminated, the growth of the unstable modes to a finite amplitude may have serious effects on the

Manuscript approved February 26, 1988.

unperturbed trajectories. It has been known that stochastic behavior[11] is an intrinsic property of perturbed Hamiltonian systems[12]. Accordingly the electron motion in an FEL will become chaotic when the sideband amplitude exceeds a certain threshold. This, in turn, will result into electron detrapping and potential loss of amplification for the electromagnetic fields.

In the present work we investigate the nonlinear effects caused by sidebands. The threshold for stochasticity, above which unbound chaotic motion occurs is determined. Once the stochastic transition takes place, the action J, a constant of motion in the unperurbed system, changes in a random manner. The ensamble average $\langle \Delta X^2 \rangle$ of any physical quantity X is described by a diffusion equation. Diffusion of the action invariant provides a measure of the leakage rate across the separatrix. If D is the effective diffusion coefficient in action space, then the diffusion length $L_d = J_s^2/D$ signifies the length over which approximately half of the deeply trapped particles get detrapped, where J_s is the action at the separatrix. We show that a single frequency sideband at one tenth (1/10) or less of the carrier amplitude suffices to spread irregular motion over most of the trapped particle domain. However, given that the interaction time of an electron in a FEL is short, we are mainly concerned on **how fast** this diffusion occurs. The diffusion rate increases and the diffusion length L_d decreases with increasing sideband amplitude(s). Thus a critical sideband level a_c can be defined above which the diffusion length L_d becomes shorter than the wiggler length L_w. Obviously the power level for the sidebands in a FEL must not exceed a_c, otherwise extensive diffusion and premature detrapping will occur. On the other hand enough electrons can remain trapped during the interaction period even though their motion has turned stochastic when the critical amplitude a_c is much larger than the threshold for stochasticity α_s.

A clear-cut relation between the diffusion rate under constant total sideband power and the type of the excited sideband spectrum is discovered. More specifically we observe three regimes in the simulation parameters defining the sideband spectrum, corresponding to a narrow, a wide discrete and a wide continuous spectrum. The transition from one spectral type to another is accompanied by an abrupt change in the diffusion rates. In all cases we find the diffusion coefficient proportional to the ratio of the total power in the sidebands to the FEL carrier power. The coefficients of this proportionality depend on the spectral type. A general conclusion is that the diffusion rate, under constant sideband over carrier power ratio, decreases with increasing number of spectral components. The diffusion rate for a single sideband frequency exceeds that of a broad continuous spectrum by orders of magnitude while a broad discrete spectrum causes intermediate diffusion rates.

The analysis also shows that the normalized diffusion coefficient does not depend on the beam energy γ_r. The numerical results agree well with the theory.

II. TRANSITION TO STOCHASTICITY

We consider relativistic electrons streaming along the z-direction through the periodic magnetic wiggler and the radiation fields of the carrier and the sideband. We assume a circularly polarized static wiggler and constant amplitude electromagnetic waves. We assume that all waves propagate with the speed of light c, ignoring the small correction of order ω_p^2/ω_r^2 from the dielectric contribution of the beam. Electrostatic contributions to the fields are neglected for operation in the Compton regime. Both the wavenumber $k_w(z)$ and the amplitude $A_w(z)$ of the wiggler may change slowly in z on a scale length much longer than the wiggler wavelength $2\pi/k_w$. The FEL main signal wavenumber k_r is doubly Doppler upshifted from the wiggler wavenumber k_w,

$$k_r = 2\gamma_z^2 k_w \ , \tag{1}$$

with the upshifting factor $\gamma_z = (1 - \beta_r^2)^{-1/2}$ and $\beta_r = \omega_r/c(k_r+k_w)$.

We normalize the time t to ω_r^{-1}, the length z to k_r^{-1}, the mass to m_e and the vector potentials according to $a_i = |e|A_i/m_e c^2$ where the subscripts w, r, and s stand for wiggler, carrier and sideband respectively. The dimensionless Hamiltonian describing the electron motion for small deviations $\tilde{\gamma} = \gamma - \gamma_r$ from the resonant energy γ_r is given by,

$$H(\tilde{\gamma},\psi;z) = \frac{k_w}{\gamma_r}\tilde{\gamma}^2 + \frac{a_w a_r}{\gamma_r}(\cos\psi + \psi\sin\psi_r) + \frac{a_w a_s}{\gamma_r}\cos(\psi - \delta_s z), \tag{2}$$

where ψ_r and δ_s are defined by,

$$\frac{d}{dz}\gamma_r = -\frac{k_r a_w a_r}{\gamma_r}\sin\psi_r \ , \qquad \delta_s = \frac{k_w}{k_r}\left(k_s - k_r\right), \tag{3}$$

The electron motion is parametrized by the traveled length z inside the wiggler. The term $\sin\psi_r$ parametrizes the rate of change for the resonant energy caused by the change in the wiggler wavelength where $\psi_r=0$ corresponds to an untapered wiggler. The term δ_s in the sideband phase is the Doppler downshifted difference between the signal and the sideband wavenumbers.

In the absence of sidebands, $a_s=0$, the Hamiltonian H_o is integrable. The unperturbed trajectories in the ponderomotive well are given by $H_o(\tilde{\gamma},\psi) = K$ where K is the reduced energy in the ponderomotive frame. These orbits take the simplest possible form expressed in terms of the action-angle variables (J, θ), defined as,

$$J = \frac{1}{2\pi}\oint d\psi\ \tilde{\gamma}(K,\psi), \qquad \theta = \frac{\partial}{\partial J}\int^{\psi} d\psi'\,\tilde{\gamma}(K,\psi')\ , \tag{4}$$

where $K = H_o(J)$ and the path of integration is over the unperturbed orbits. Hamiltonian (2) is transformed under the canonical transformation defined by Eq. (4) into,

$$H(J,\theta;z) = H_o(J) + \frac{a_w a_s}{\gamma_r}\sum_{n=0}^{\infty} Q_n^{+}(J)\cos(n\theta + \delta_s z) + Q_n^{-}(J)\cos(n\theta - \delta_s z), \tag{5}$$

$Q_n^{+}(J)$ are the Fourier coefficients obtained by the decomposition of the perturbing sideband phase $\psi(J,\theta) - \delta_s z$ into harmonics of the angle θ, where $\psi(J, \theta)$ is obtained by inverting Eq. (4). In case of constant parameter wiggler J, θ and $Q_n(J)$ are expressed in closed forms given in Ref. 11.

$H_o(J)$ is independent of θ so the unperturbed orbits in (J, θ) space are straight lines J = constant, $\theta = \theta_o + \kappa_b(J)z$. The synchrotron wavenumber $\kappa_b(J)$ is connected to the bounce length L_b and the synchrotron frequency in the laboratory frame $\omega_b(J)$ with the relation

$$\kappa_b(J) = \frac{dH_o(J)}{dJ} = \frac{2\pi}{L_b(J)} = \frac{\omega_b(J)}{c\beta_z} . \tag{6}$$

According to Eq. (5) new resonances emerge when a sideband is turned on. Defining the phase of the nth sideband induced harmonic $\theta^{(n)} = n\theta \pm \delta_s z$, the stationary phase condition reads,

$$\pm n\kappa_b(J) - \delta_s = 0, \qquad \text{or} \qquad \pm n\beta_z\omega_b(J) - \delta_s = 0. \tag{7}$$

Thus particles, originally in unperturbed orbits $J = J_n$, resonate with the sideband when the nth harmonic of their synchrotron period $\omega_b(J_n)$ matches the downshifted frequency difference between the sideband and the carrier signal. When the sideband amplitude becomes finite these resonant particles get trapped inside secondary ponderomotive wells corresponding to the different harmonics in Hamiltonian (4). Thus a single frequency sideband causes chains of secondary islands to appear inside the original primary island. The structure of the phase space is shown in Figs. 1a – 1d. They are surfaces of section, created by numerically integrating the original equations of motion from Hamiltonian Eq. (2) and then recording the intersection point of each trajectory with the plane $z = 2\pi/\delta_s$. The γ, ψ coordinates are plotted on the left side in Figs. 1a-1d. The plots on the right side show the same surfaces of section in action-angle variables, produced by the transformations (4).

The presence of just one sideband frequency suffices to transform the regular coherent motion, such as in Figs. 1a and 1c, to the irregular unbounded motion shown in Figs. 1b and 1d when the sideband amplitude exceeds a certain amplitude α_s regarded as the stochasticity threshold. The mechanism for this radical change in behavior can be briefly described as follows. The trajectories emanating from the unstable fixed points (X-points) of a secondary island corresponding to a given harmonic n do not actually join smoothly around that island. They intersect infinite times with each other[11,12] due to the effect of the other harmonics $n' \neq n$, creating a thin layer of fussy motion that surrounds each island chain. As the amplitude a_s increases the width of each island increases and so does the thickness of the stochastic layer around that island. At a given point the stochastic layers around the two neighboring island chains n and n+1 overlap[13], allowing particles to hop from one island to another. This signifies the beginning of unbounded, random motion in J characterized as stochastic diffusion.

The overlapping condition that defines the transition to stochasticity is usually expressed by[11,13]

$$\delta J_n + \delta J_{n+1} \geq \frac{2}{3} \Delta J_n \ , \tag{8}$$

where δJ_n, δJ_{n+1} are the separatrix half-widths and $\Delta J_n = J_{n+1} - J_n$ is the distance between the separatrix centers for the n and n+1 harmonics respectively. δJ_n and ΔJ_n are evaluated from Hamiltonian (5) using the familiar methods of nonlinear dynamics in terms of the sideband amplitude a_s and the shear $d\kappa_b(J)/dJ$. The detailed computation, performed in Ref. 14, shows that the threshold for stochasticity α_s depends primarily on the frequency separation $(\omega_s - \omega_r)/\omega_r$ and less on the other FEL parameters. The ratio α_s/a_r varies in a semioscillatory manner between typical values of 0.1 and 1. The above threshold was derived for a single sideband frequency. Multifrequency sideband spectra generally tend to reduce α_s.

III. NARROW SPECTRUM DIFFUSION

We examine first the diffusion caused by a narrow sideband spectrum, meaning that the width of the sideband spectral line is much less than its distance from the carrier frequency. We model this case by a single frequency sideband. Figures 1b and 1d show a typical phase portrait when the sideband amplitude a_s is slightly above α_s. Two different kinds of regions coexist: a stochastic regime where diffusive behavior prevails, interrupted here and there by islands of regular motion, remnants of the original regular motion. The stochastic regimes are interconnected allowing particle transport, hence α_s was defined as global stochasticity threshold.

Figure 2 illustrates the stochastic transition by plotting selected particle trajectories J(z) with z measured in wiggler periods. The sideband amplitude increases from well below α_s in Fig. 2a to well above α_s in Fig. 2c. The initial conditions for the particles remain the same in all cases. The oscillations in Fig. 2a have the period of a particle trapped around a secondary island.

When a_s is increased well above α_s the chaotic motion engulfs almost 100% of the phase space. Total stochastization of the island interior occurs roughly when the sideband amplitude grows to the point where the stable fixed point $\tilde{\gamma} = 0$, $\psi = \pi + \psi_r$ at the centre of the original island becomes unstable.

The two questions of practical interest are: (a) what percentage of the particles will eventually get detrapped and (b) how fast do they leak outside the separatrix. For our uniform initial distribution the maximum fraction of particles becoming detrapped equals the fraction of the inside the separatrix area that becomes chaotic. In Fig. 3a we plot the fraction f_d of the particles that cross the original separatrix J_s as a function of the traveled wiggler length for values of a_s/a_r below the threshold for extended stochasticity. In all cases an initial stage of quick diffusion is followed by a long period where the average number of untrapped particles remains practically constant. The results are consistent with the existence of a boundary in phase space (KAM surface) separating two regimes: the one of unbounded, chaotic motion from the one filled with regular, coherent orbits of particles that remain trapped. Only electrons in the area between the last integrable surface and the old unperturbed separatrix will diffuse until that area is depleted. A fraction $1 - f_d$ of the original primary island area will remain trapped for an arbitrarily long time, as long as a_s remains below the threshold α_s associated with the particular sideband frequency. This fraction is shrinking as a_s increases and the bucket "peels off". The situation when a_s exceeds α_s is shown in Fig. 3b. The fraction of untrapped particles f_d increases monotonically, reaching 1 in some cases, meaning complete absence of particle confinement in the bucket. All particles can eventually escape with a rate that increases with increasing a_s. The length over which approximately half of the initially trapped particles get detrapped will be discussed in the next section, in comparison with the diffusion rates from other types of sideband spectra. It is evident from the results in Fig. 3 that the diffusion length at constant amplitude a_s depends on the sideband frequency.

IV. BROAD FREQUENCY BAND DIFFUSION

So far stochastic electron detrapping caused by a single frequency sideband has been examined. Here, we consider the situation when many sideband frequencies have been excited. We will make a distinction between a continuous and a discrete spectrum. In case of a discrete spectrum the distance between two nearby sideband frequencies is much larger than the width of an individual spectral line. In the opposite case when various peaks in the spectrum merge together we will talk about a continuous spectrum. We may model numerically both cases by introducing a modulation in the sideband phase of Hamiltonian Eq. (4),

$$H(\tilde{\gamma},\psi;z) = \frac{k_w}{\gamma_r}\tilde{\gamma}^2 + \frac{a_w a_r}{\gamma_r}(\cos\psi + \psi \sin\psi_r) + \frac{a_w a_s}{\gamma_r}\cos(\psi + A \sin \nu z - \delta_s z), \tag{9}$$

that is transformed in action-angle variables as,

$$H(J,\theta;z) = H_o(J) + \frac{a_w a_s}{\gamma_r}\sum_{m=-\infty}^{\infty} J_m(A) \sum_{n=0}^{\infty} Q_n^+(J) \cos[n\theta + \delta_s(m)z] + Q_n^-(J) \cos[n\theta - \delta_s(m)z] \,. \tag{10}$$

The frequency mismatch values $\delta_s(m)$ and the corresponding sideband frequencies $\omega_s(m)$ are given by,

$$\delta_s(m) = \delta_o + m\,\nu, \qquad \omega_s(m) = \omega_{so} + 2\,m\,\gamma_z^2\nu, \tag{11}$$

where $\omega_{so} = \omega_r + 2\gamma_z^2\delta_o$. Since the Bessel function coefficients become vanishingly small, $J_m(A) << 1$ for $A > m$, the width of the spectrum is given by $D\delta_s \sim A\,\nu$ or $D\omega_s \simeq 2\gamma_z^2 A\nu$.

In order to examine the connection between diffusion rates and the types of the sideband spectra, we divide the latter into three general categories: narrow, broad discrete and broad continuous. The passage from one regime to the other is characterized by abrupt changes in the diffusion coefficients. Both cases of the broad spectrum are characterized by a width $D\omega_s$ in the excited frequencies that is larger than the upshifted synchrotron frequency ω_b, $D\omega_s > 2\gamma_z{}^2\omega_b$, or equivalently,

$$\nu > \frac{\omega_b}{A} , \tag{12}$$

with $A >> 1$. The further distinction between discrete or continuous spectrum is related to the separation between nearby frequencies. We find that when $\omega_b/A^{1/2} > \nu > \omega_b/A$ the diffusion rate agrees well with the quasilinear diffusion coefficient. A different coefficient is derived for the case when $\nu > \omega_b/A^{1/2} > \omega_b/A$, in agreement with the numerical simulations. Consequently, the separation $\omega_b/A^{1/2}$ between nearby modes marks the transition from a discrete to a continuous type of behavior.

(a). Broad Discrete Spectrum.

We now evaluate the diffusion coefficient for a broad, discrete spectrum associated with frequency separation $\nu > \omega_b/A^{1/2}$. Due to the presence of many frequencies in the spectrum J(z) executes a complicated oscillatory motion with the average $< J >$ changing very little most of the time. J however receives a large kick ΔJ near resonances, where the phase $\Phi = \psi + A \sin \nu z \pm \delta_s z$ of the multifrequency Hamiltonian (9) varies slowly. The resonant condition is

$$\frac{d\Phi}{dz} = \frac{k_w \tilde{\gamma}}{\gamma_r} + A\, \nu \sin \nu z_i \pm \delta_s = 0, \tag{13}$$

at some $z = z_i$. Given that $k_w \tilde{\gamma}/\gamma_r \leq \omega_b$, collective effects due to many frequencies are important for the resonance in Eq. (13) when $A\, \nu > \omega_b$. On this basis inequality (12) signifies the transition from a narrow to a broad spectrum. Let us consider the case $A\, \nu >> \omega_b$. Then the resonances occur at $z_i \simeq i\pi/\nu$, i integer, and the interval between successive resonances is $\Delta z \simeq \pi/\nu$. Expanding the phase $\Phi(z)$ in the equation of motion for J around the resonance z_i, and extending the limits of the z integration from $-\infty$ to $+\infty$ we obtain, for $A\, \nu^2 >> \omega_b^{\,2}$,

$$DJ_i \simeq \frac{a_w a_r}{\gamma_r} \left(\frac{2\pi}{A\, \nu^2}\right)^{1/2} \cos\left(\Phi_n \pm \frac{\pi}{4}\right) \left|V(\psi_{mx}) - V(\psi_i)\right|^{1/2}, \tag{14}$$

where $V(\psi) = \cos\psi + \psi \sin\psi_r$. The resonant phases Φ_i between two successive jumps ΔJ become quickly decorrelated when a_s grows above the stochasticity threshold. Assuming complete decorrelation in Φ_i between two successive jumps we obtain,

$$D_w = \frac{2\, \langle \Delta J^2 \rangle}{\Delta z} = \frac{2\, a_w^{\,2} a_s^{\,2}}{\gamma_r^{\,2} A\, \nu}, \tag{15}$$

where <....> is the ensemble average over ψ_i, Φ_i. For practical purposes it is more convenient to rescale the diffusion coefficient so that the distance $\underline{z} = z/\lambda_w$ is measured in terms of wiggler wavelengths and the action $\underline{J} = J/J_s$ signifies the location relative to the separatrix. In these units, using Eq. (4) for J_s and setting $\nu / \omega_b(0) = r$ we obtain

$$\underline{D}_w = \frac{2\pi}{k_w} \frac{D_w}{J_s^2} = \frac{\pi^3}{8} \frac{g^2}{A\ r} \left(\frac{a_w a_r}{1 + a_w^2}\right)^{1/2} \frac{a_s^2}{a_r^2} . \qquad (16)$$

The term g is a scaling factor, the ratio of the untapered wiggler separatrix to that of a tapered wiggler, depending only on ψ_r, $g = J_s(\psi_r{=}0)/J_s(\psi_r)$. The typical diffusion length L_d, the traveled distance inside the wiggler over which the average trapped particle crosses the separatrix is estimated from $\langle\Delta J^2\rangle = J_s^2 = D_w L_d$. Thus, the diffusion length in wiggler periods $N_d = L_d/\lambda_w$ is the inverse of $\underline{D}_w$, $N_d = 1/\underline{D}_w$.

(b). Broad Continuous Spectrum

Next we consider the case of a sideband wave package,

$$a_s(z,t) = \frac{1}{2\pi} \int dk_s a_s(k_s)\ e^{ik_s z - i\omega(k_s)t} , \qquad (17)$$

of finite spectral width Dk_s centered around k_{so}. Our purpose is to obtain the diffusion coefficient for a continuous spectrum using the methods of the quasilinear theory.

One condition for the applicability of the quasilinear theory is that the phase mixing due to Dk_s occurs much faster than the bounce period around a secondary island in phase space. This way electrons, that otherwise would execute periodic orbits around some fixed point, lose coherency fast enough and a random motion of the Fokker-Planck type sets in. This condition is satisfied for spectral widths Dk_s of order,

$$\frac{Dk_s}{k_s} >> 2\ \gamma_z^2\ \frac{\kappa_b^2(0)}{k_r^2}\ . \tag{18}$$

The overlapping condition among nearby island chains, similar to Eq. (8), must also be satisfied. The threshold for overlapping is much lower than α_s related to the single sideband frequency. Then the evolution of the initial distribution $f_o(J)$ is globally described by a diffusion equation,

$$\frac{\partial f}{\partial z} = \frac{\partial}{\partial J} D_q(J) \frac{\partial f}{\partial J}\ .$$

Applying the standard procedures of the quasilinear theory[15] and taking the limit of small growth rate for the sidebands, $Im(k_s)/k_s << 1$, we obtain the quasilinear diffusion coefficient,

$$D_q(J) = \frac{k_r\ a_w^2\ \gamma_z^2\ \underline{W}_s(k_{so})\ k_w\ J}{\gamma_r^3 \omega_b(J)}\ . \tag{19}$$

We express $D_q(J)$ in normalized units, with the wiggler wavelength λ_w as the unit length and the action J_s at the separatrix as the unit action. Chosing the value $J = J_s/2$ inside $D_q(J)$ we obtain an estimate for the effective diffusion coefficient associated with the uniformly filled distribution,

$$\underline{D}_q = \frac{2\pi}{k_w} \frac{D_q(J)}{J_s^2} \sim \frac{\pi}{4} \frac{g\, a_w a_r}{\Lambda\, (1+a_w^2)} \frac{W_s(t)}{Dk_s\, W_r(t)} , \qquad (20)$$

where g is the same scaling factor as in Eq. (16).

Note that both expressions (16) and (20) for D, corresponding to the two different spectral types, are independent of γ_r. Thus, for the same wiggler parameters and total sideband power, the detrapping distance in wiggler periods is independent of the electron beam energy.

A more detailed treatement of the subject, including the derivation of the formulas (16) and (20) is found in Ref. 14.

V. NUMERICAL RESULTS

The trajectories J(z) on the left side of 4a-4c are generated by the same initial conditions for the electrons, the same FEL parameters a_w, a_r and k_w, and the same averaged sideband power $\langle a_s^2 \rangle$. The spectral parameters A and ν however are different so that each of the cases (a) (b) and (c) corresponds to one of the three spectral types defined earlier. The dashed line marks the position of the unperturbed separatrix J_s. Without the sidebands the trajectories would be straight lines. The corresponding distribution functions f(J,z) are plotted on the right-hand side at the beginning, $z = 0$, halfway inside, $z = 50\lambda_w$, and at the end, $z = 100\lambda_w$ of the wiggler.

The numerically computed diffusion length in wiggler periods $N_d = 1/\underline{D}$ is plotted against the sideband to carrier power ratio $P = \sum_n a_s^2(\omega_n)/a_r^2 = W_s/W_r$ in Fig. 5 for the three different types of spectra . We have integrated numerically the equations of motion for 400 particles of a uniform initial distribution inside the bucket. The field intensities remained constant at $a_r = 5\times10^{-5}$, $a_w = 2$ and $\gamma_r = 25$. A clear separation in the diffusion rates is observed among the various spectral types. The narrow frequency results (triangles) were obtained using the Hamiltonian (2) with a single sideband frequency $\omega_s/\omega_r = 1.016$. The results for a broad discrete spectrum (circles) were obtained using (9) with $A = 20$, $\omega_s/\omega_r = 1.016$ and $\nu = 0.5\ \delta_s$. The continuous spectrum (squares) was modeled by $A = 100$, $\nu = 0.05\ \delta_s$. The solid lines correspond to the theoretical results of Eqs. (16) and (20), in good agreement with the numerics. Theoretical predictions for the single frequency case were not made. We stress however the difference between single frequency results and quasilinear theory in this case. The agreement that has been observed in some other cases is not generic but particular to certain systems.

Finally, in Fig. 6 we plot the diffusion coefficient for a uniformly filled bucket as a function of the energy γ_r, fixing the wiggler parameters. It is clear that the diffusion rate (measured again in number of wiggler periods) is independent of the beam energy, provided the synchrotron frequency ω_b stays in the same parameter regime.

Once the diffusion coefficients are known some estimate can be made of the related reduction in efficiency over the wiggler length. Assuming that $n_b(z) = n_b(0) \exp(-z/L_d)$ the number of detrapped particles between z and z + Δz is $\Delta n(z) = n_b(0)L_d^{-1} \exp(-z/L_d)\ \Delta z$. These particles gave up an amount of energy $\Delta E(z) = [\gamma_r(0) - \gamma_r(z)]\ \Delta n(z)$ as radiation. Integrating $\Delta E(z)$ over the wiggler length for a linearly tapered wiggler $\gamma_r(z) = \gamma_r(0) - \Delta\gamma\ z / L_w$ we find the total energy extracted from the particles that were detrapped at some point inside the wiggler. Adding the contribution $[\gamma_r(0) - \gamma_r(L_w)]\ n_b(L_w)$ from the particles that remained trapped throughout the wiggler length we come up with,

$$\eta(L_w) = \eta_o \frac{L_d}{L_w} \left(1 - \exp\left(- \frac{L_w}{L_d}\right) \right), \tag{21}$$

where $\eta_o = \Delta\gamma/\gamma_r(0)$ is the efficiency without induced diffusion. The loss of amplification will in general be distributed among all the radiation modes and (21) reflects the total power loss in all frequencies.

Acknowledgement

This work supported by SDIO and managed by SDC.

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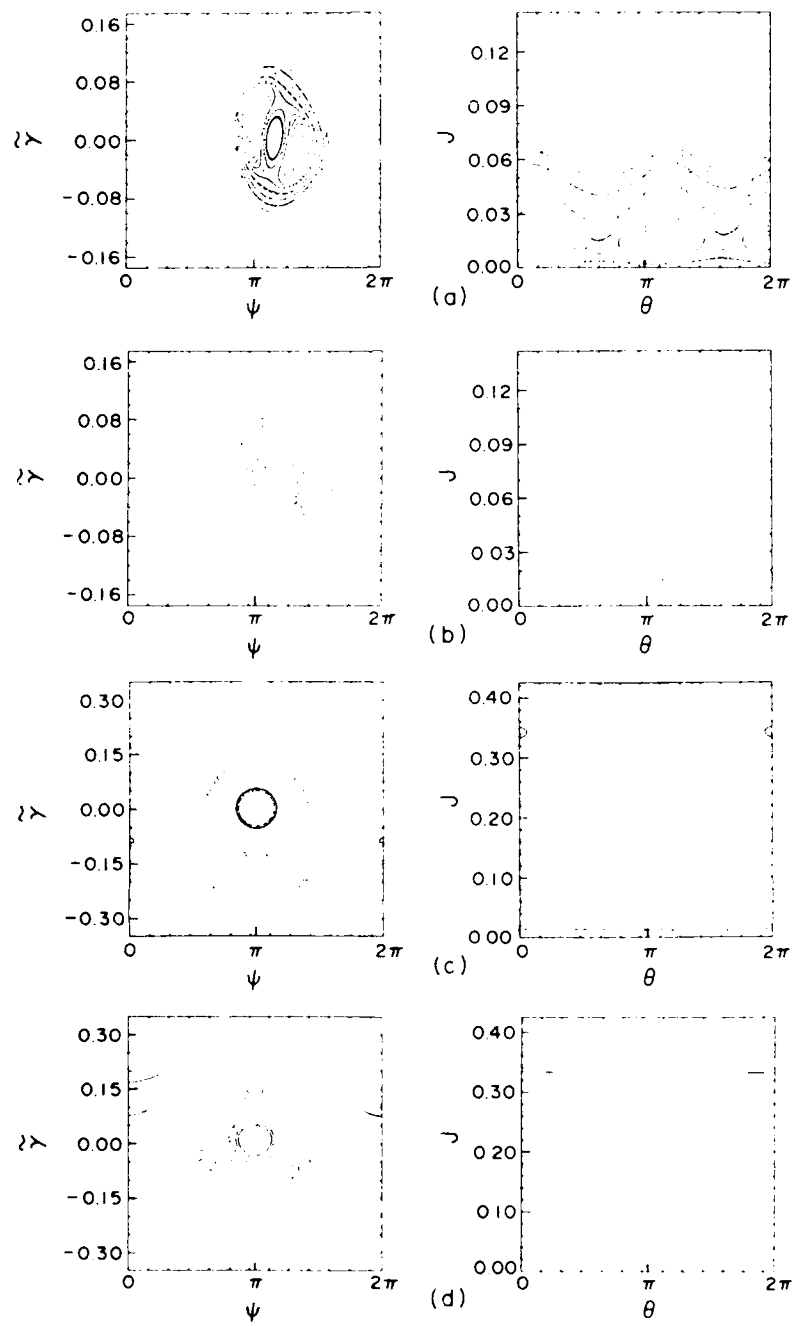

Figure 1. Plots of surfaces of section in both $\tilde{\gamma}$, ψ and J, θ representations, illustrating transition to chaotic behavior. The parameters are $a_w = 2$, $a_r = 5\times10^{-5}$ and initial $\gamma_r = 25$. Figures (a) and (b) correspond to tapered wiggler, $\psi_r = 4\pi/3$, with $\omega_s/\omega_r = 1.016$ and (a) $a_s = 2\times10^{-6}$, (b) $a_s = 1\times10^{-5}$. Figures (c) and (d) correspond to untapered wiggler with $\omega_s/\omega_r = 1.024$ and (c) $a_s = 2\times10^{-6}$, (d) $a_s = 1\times10^{-5}$.

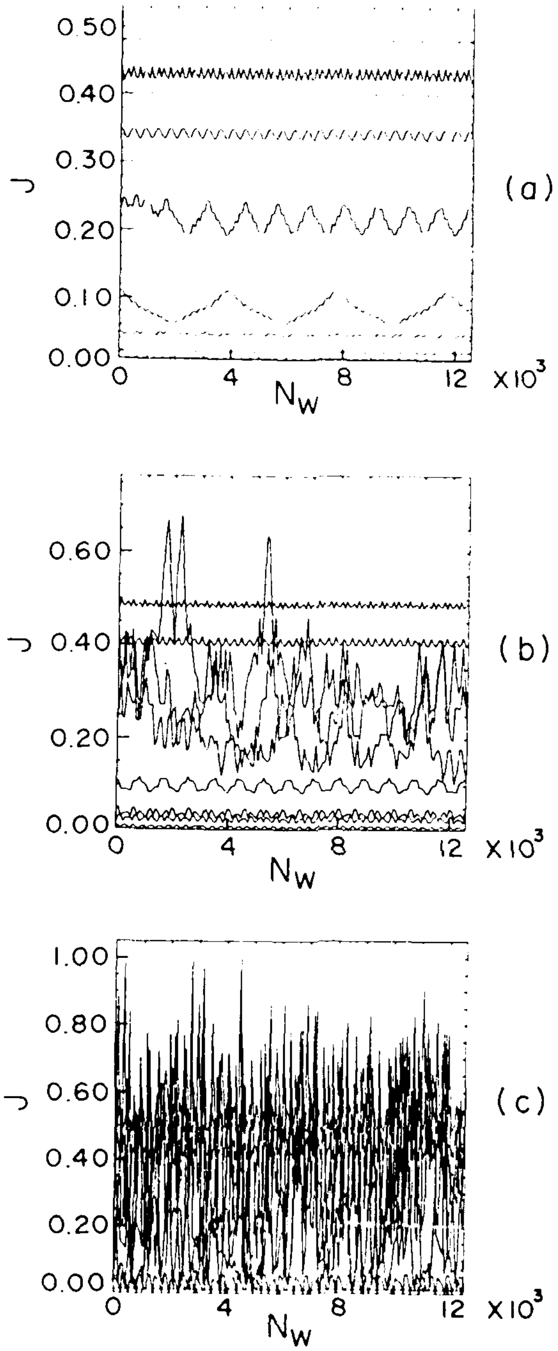

Figure 2. Transition to chaotic behavior. The action J is plotted against the number of wiggler periods N_w. The parameters are $a_w = 2$, $a_r = 5\times10^{-5}$, $\gamma_r = 25$, $\omega_s/\omega_r = 1.024$ and (a) $a_s = 2\times10^{-6}$, (b) $a_s = 1\times10^{-5}$ and (c) $a_s = 5\times10^{-5}$.

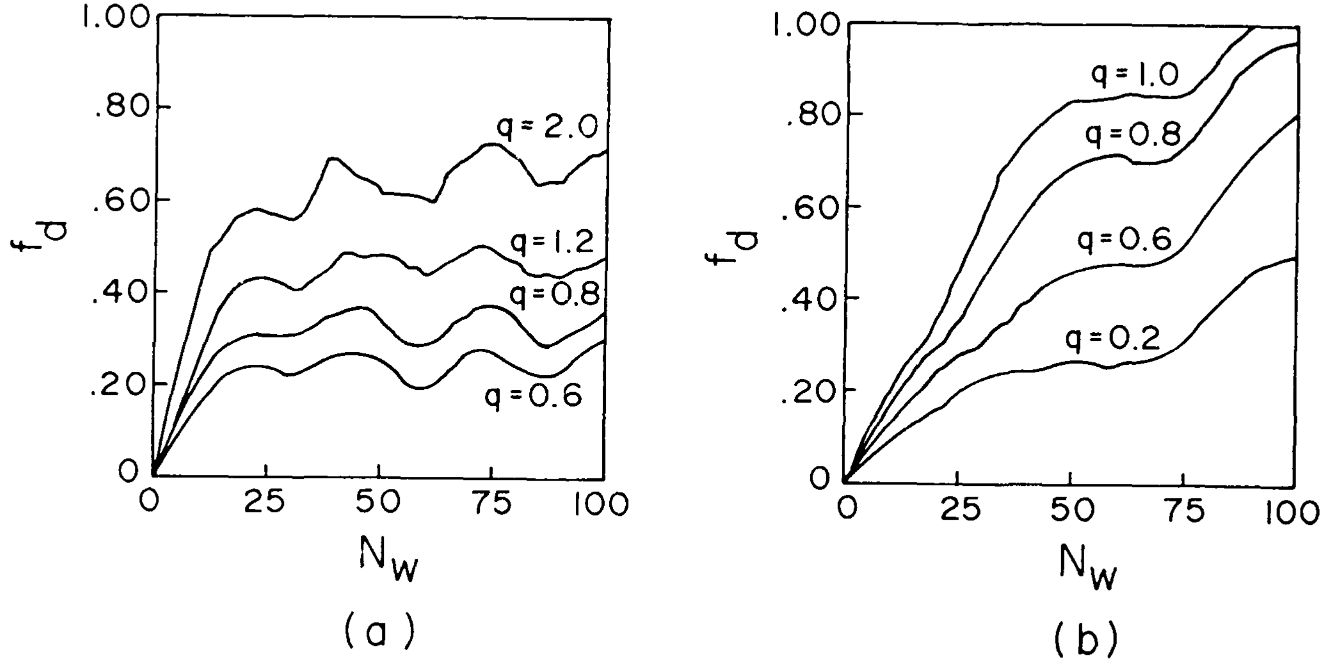

Figure 3. The fraction f_d of detrapped particles against the distance in wiggler periods $N_w = z / \lambda_w$. Different curves correspond to various sideband to carrier amplitude ratios $q = a_s/a_r$. Parameters are $a_w = 2$, $a_r = 5\times10^{-5}$, $\gamma_r = 25$ and (a) $\omega_s/\omega_r = 1.016$, (b) $\omega_s/\omega_r = 1.024$.

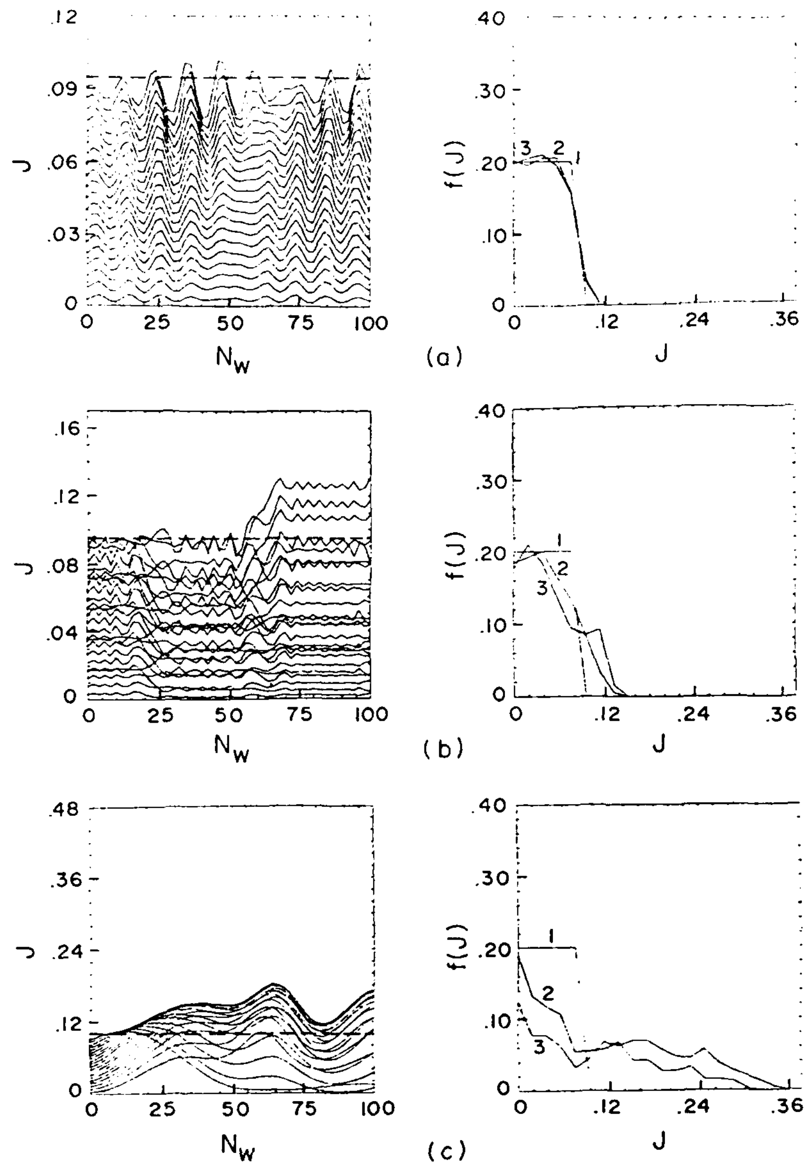

Figure 4. Particle response under different sideband spectra. On the left-hand side we plot the action J against the number of wiggler periods N_w for selected particles. On the right-hand side we plot the corresponding distribution function f(J) at N_w = 0, 50, and 100 . In all cases the total sideband fractional power $P = W_s/W_r = 0.36$ and $a_w = 2$, $a_r = 5\times10^{-5}$, $\gamma_r = 25$. (a) corresponds to a discrete sideband spectrum with $\omega_s/\omega_r = 1.016$. (b) corresponds to a wide discrete spectrum and (c) to a wide continuous spectrum.

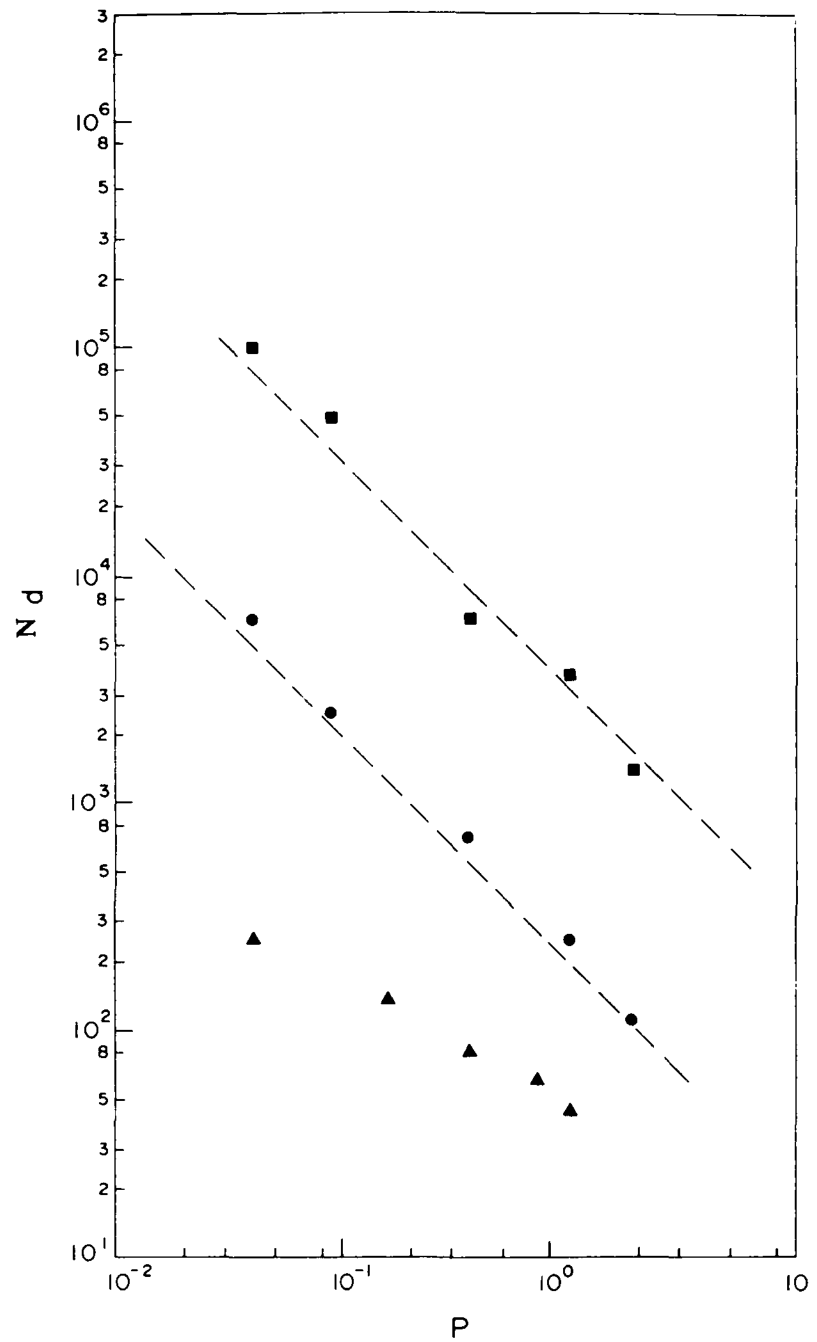

Figure 5. The e-folding diffusion length $N_d = 1/D$ as a function of the sideband fractional power $P = W_s/W_r$ for $a_w = 2$, $a_r = 5\times10^{-5}$, $\gamma_r = 25$ and $\omega_s/\omega_r = 1.016$. Squares correspond to a continuous, dots to a wide discrete and triangles to a single frequency spectrum. The upper and lower solid lines correspond to the theoretical results Eqs. (16) and (20) respectively.

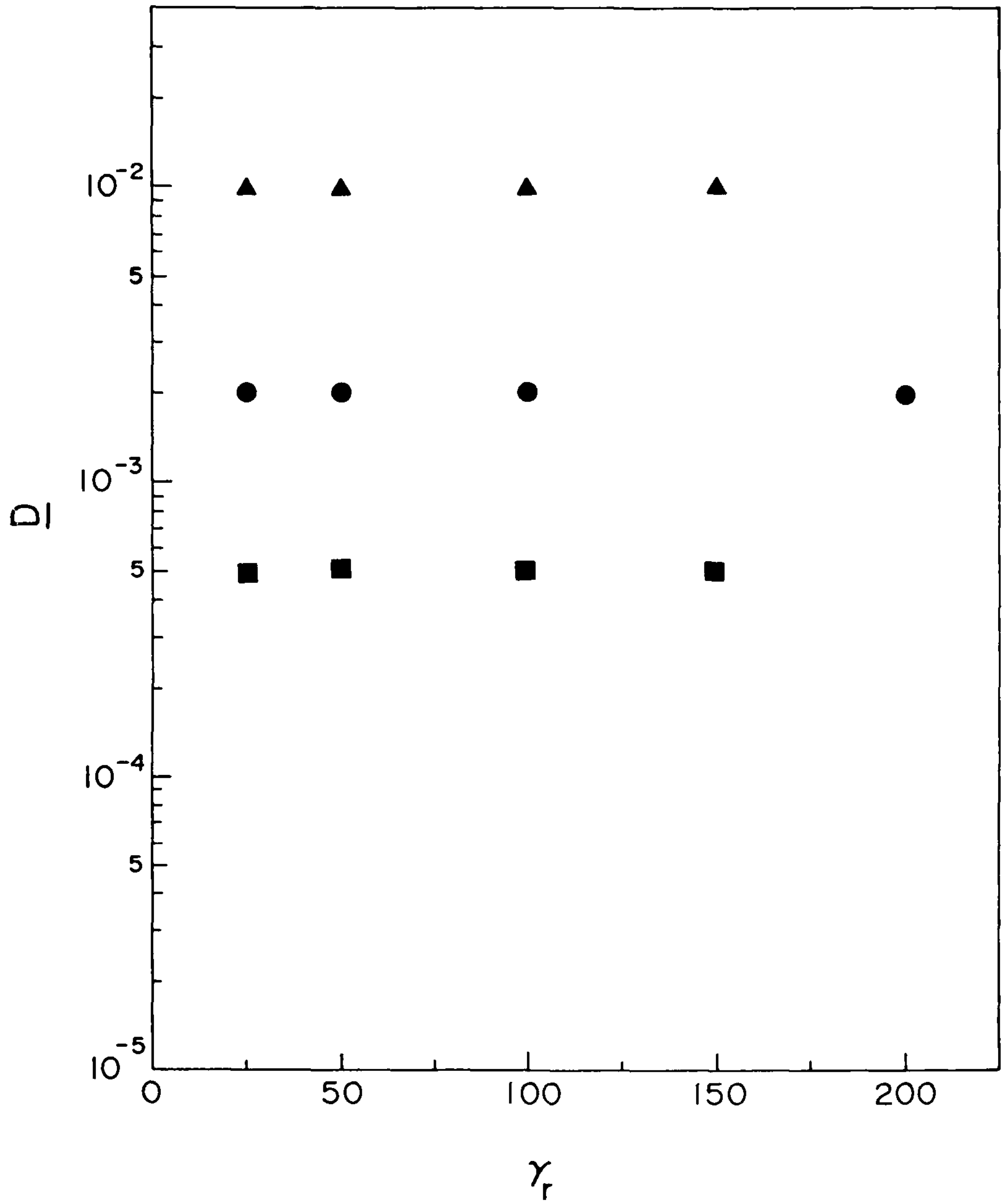

Figure 6. Plots of the normalized diffusion coefficient $\underline{D}$ against the electron energy γ_r for $a_w = 2$, $a_r = 2\times10^{-4}$ and $a_s = 7.5\times10^{-5}$. Squares correspond to a continuous, dots to a wide discrete and triangles to a single frequency spectrum.

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