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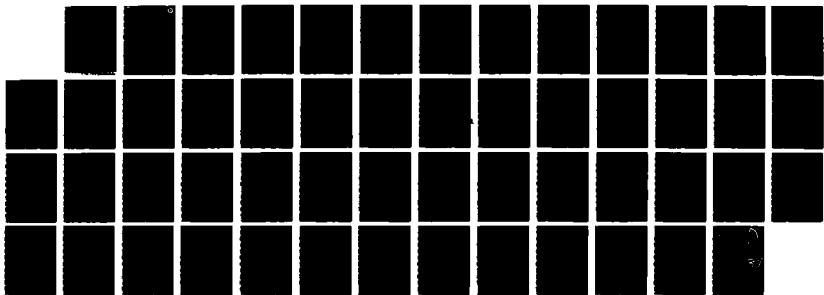
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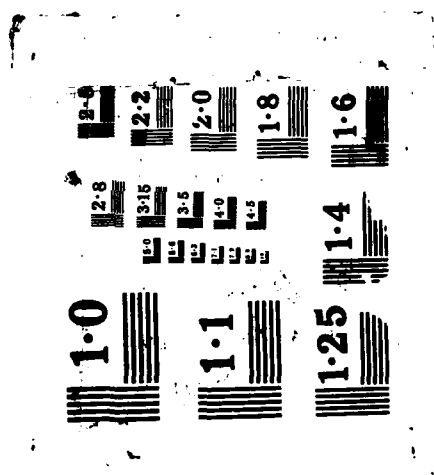
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The Free Electron Laser Sideband Instability Reconsidered

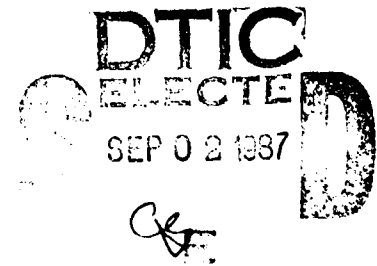
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August 12, 1987



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SECURITY CLASSIFICATION OF THIS PAGE

REPORT DOCUMENTATION PAGE

1a. REPORT SECURITY CLASSIFICATION UNCLASSIFIED			1b. RESTRICTIVE MARKINGS		
2a. SECURITY CLASSIFICATION AUTHORITY			3. DISTRIBUTION/AVAILABILITY OF REPORT Approved for public release; distribution unlimited.		
2b. DECLASSIFICATION/DOWNGRADING SCHEDULE					
4. PERFORMING ORGANIZATION REPORT NUMBER(S) NRL Memorandum Report 6028			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			
6c. ADDRESS (City, State, and ZIP Code) Washington, DC 20375-5000		7b. ADDRESS (City, State, and ZIP Code)			
8a. NAME OF FUNDING/SPONSORING ORGANIZATION US Army Strategic Defense Communication	8b. OFFICE SYMBOL (If applicable)	9. PROCUREMENT INSTRUMENT IDENTIFICATION NUMBER			
8c. ADDRESS (City, State, and ZIP Code) Huntsville, AL 35807-3801	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) The Free Electron Laser Sideband Instability Reconsidered					
12. PERSONAL AUTHOR(S) Riyopoulos*, S. and Tang, C.M.					
13a. TYPE OF REPORT Interim	13b. TIME COVERED FROM TO	14. DATE OF REPORT (Year, Month, Day) 1987 August 12		15. PAGE COUNT 53	
16. SUPPLEMENTARY NOTATION *Science Applications Intl. Corp., McLean VA.					
17. COSATI CODES			18. SUBJECT TERMS (Continue on reverse if necessary and identify by block number)		
FIELD	GROUP	SUB-GROUP	FEL instabilities Sidebands		
			Canonical formalism for kinetic equation		
19. ABSTRACT (Continue on reverse if necessary and identify by block number) The problem of sideband growth in a free electron laser (FEL) is studied using canonical formalism. The exact unperturbed orbits, without sidebands, are obtained in action-angle variables for all particles. Integration of the linearized Vlasov equation with perturbing sidebands over the unperturbed orbits yields the sideband growth for both trapped and untrapped particles. The unperturbed distribution is in an equilibrium with the main signal field. It is found that upper and lower sidebands that are symmetric relative to the FEL frequency have opposite growth rates. There is no differentiation in the magnitude of the growth between upper and lower sidebands. The stability is determined by the sign of $df_0/d\omega_b$, i.e., the relative population of oscillation quanta $\hbar\omega_b$, ω_b = bounce frequency around resonance. The shear $d\omega_b/dJ$, where J is the action variable, is stabilizing and distributions with gradients df/dJ localized near the separatrix have the minimum growth rates. The structure and scaling of the unstable spectrum are different from previous results obtained for electrons localized at the bottom of the ponderomotive well.					
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) 767-4148		22c. OFFICE SYMBOL Code 4790

DD FORM 1473, 84 MAR

83 APR edition may be used until exhausted
All other editions are obsolete.

SECURITY CLASSIFICATION OF THIS PAGE

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THE FREE ELECTRON LASER SIDEBAND INSTABILITY RECONSIDERED

I. INTRODUCTION AND SUMMARY

The growth of parasitic modes at frequencies near the main signal frequency during high power FEL operation was theoretically predicted^{1,2} in the early 1980's. Since then there has been ample numerical^{3,4} and experimental^{5,6} evidence of sideband excitation in constant wiggler FELs. Unstable modes in variable wiggler FELs have also been observed in simulations⁷⁻⁹ and recently in experiment¹⁰. Sidebands degrade the main signal efficiency and optical quality by channeling a considerable fraction of the power into parasitic frequencies. The performance of the mirrors in an oscillator can be harmed from the modulation of the wave envelope caused by the sidebands. Last, but not least, interaction among nearby sidebands above a certain amplitude may lead to chaotic particle motion, loss of trapping and incoherent radiation.

The above have stimulated a considerable amount of theoretical work focused on sideband growth. Simple one-dimensional configurations that are analytically tractable have been used to model the situation. Two lines of approach have been considered. The single particle picture regards the particle trajectories as functions of the initial conditions and computes the growth by ensemble averaging over initial distributions⁷⁻⁹. The alternative approach assumes some adiabatic equilibrium between the particles and the main signal and examines the stability of the perturbations induced by the sidebands, solving the kinetic equation^{11,12}. Because of the equilibrium assumption the

kinetic method is more appropriate for FELs operating at low gain.

In both treatments so far, analytic results have been obtained only for particles localized near the bottom of the ponderomotive well. This implies the following limitations: The sideband spectrum becomes discrete

$$\omega_s = \omega_r \pm (k_r/k_w)n\omega_b(0), \quad k_r/k_w \approx 2\gamma_z^2, \quad (1)$$

where $\omega_b(0)$ is the bounce frequency at the bottom of the ponderomotive well, k_r , k_w are the radiation and wiggler wave numbers respectively and $\gamma_z = (1 - v_z^2/c^2)^{-1/2}$. The contribution from untrapped particles and trapped particles away from the bottom is neglected. The effect of the shear $d\omega_b/dJ$, where the action J parametrizes the distance from the centre of the separatrix $J=0$, is lost. Finally no predictions can be made about the saturation level of potentially unstable modes.

Here canonical formalism is introduced by expressing the unperturbed particle orbits in terms of action-angle variables. The unperturbed orbits are the fast time averaged "synchrotron" oscillations of the electrons in the potential well formed by the combined action of the wiggler and the radiation signal. The perturbed kinetic equation is solved in action space, starting from an equilibrium extending over all trapped and untrapped electrons. It is found that:

(a) the spectrum becomes continuous replacing $\omega_b(0)$ by $\omega_b(J)$ in Eq. (1). The modes located at the peaks of the unstable spectrum grow faster, emerging as the discrete spectrum that is observed in simulations.

(b) More than one groups of particles are in resonance with a given sideband frequency ω_s , through different harmonics of their bounce frequency, and contribute to the growth rate.

(c) Upper and lower sidebands located symmetrically around the main signal frequency have opposite growths (complementary stability). Therefore one mode is always unstable. There is no stable distribution $f_0(J)$ except the trivial one $df_0/dJ=0$.

(d) The shear $d\omega_b/dJ$ is stabilizing. Distributions with gradients df_0/dJ localized near the separatrix are found to have the minimum growth rates because of the high shear there. This type of distribution is relevant to FELs with tapered wigglers.

(e) The growth is proportional to $[df(J)/d\omega_b(J)]$, the relative population in oscillation quanta around resonance, in agreement with the quantum mechanical interpretation.

(f) For any smooth distribution, of finite df_0/dJ , electrons at the bottom of the well have a negligible effect on stability.

(g) Previous results, finding lower sidebands having an inherently larger growth than upper sidebands, are relevant only to the limiting case of a singular δ -function distribution $f_0(J)=\delta(J)$. This case is unrealistic because a wide, smooth initial distribution in action $f_0(J)$ corresponds to even an ideal cold beam distribution in momentum $f_0(p)=\delta(p-p_0)$.

The nonlinear saturation levels for the unstable modes and the amplitude for stochastic transition can also be derived from our formalism and will be addressed in future work.

II. COMPUTATION OF THE GROWTH

We consider a circularly polarized wiggler of constant wavelength $2\pi/k_w$ and constant amplitude A_w and circular monochromatic waves for the main signal and the sideband. The total vector potential is given by

$$A = \frac{1}{2} \left[(\mathbf{e}_x - i\mathbf{e}_y) A_w e^{ik_w z} - (\mathbf{e}_x + i\mathbf{e}_y) A_r e^{i(k_r z - \omega_r t)} - (\mathbf{e}_x + i\mathbf{e}_y) A_s e^{i(k_s z - \omega_s t)} \right] + CC \quad (2)$$

where the subscripts w , r and s stand for wiggler, main radiation and sideband respectively, A_r , A_s , are constant and $A_r \gg A_s$. Dispersive effects are generally very small, of order $(\omega_p/\omega_r)^2 \ll 1$, and all electromagnetic waves propagate with the speed of light. Corrections to the real part of the wave number $k(\omega_s) = \omega_s/c$ are therefore negligible while the imaginary part of k_s can be computed from the wave-particle energy balance equation without resorting to the solution of the dispersion relation. Non linear phase slippage effects are also omitted.

The motion of the electrons in the fields of Eq. (2) is described by the Hamiltonian

$$H(\tilde{\gamma}, \psi; z) = \frac{k_w}{\tilde{\gamma}_r} \tilde{\gamma}^2 - \frac{a_w a_r}{\tilde{\gamma}_r} (\cos \psi - \psi \sin \psi_r) - \frac{a_w a_s}{\tilde{\gamma}_r} \cos(\psi - \delta_s z) \quad (3)$$

where $\tilde{\gamma} = \gamma - \gamma_r$, time t is normalized to ω_r^{-1} , length z to k_r^{-1} , mass to m_e and $a_i = eA_i/m_e c^2$. Although $k_r = \omega_r = 1$ in the above units we write them down explicitly to avoid some confusion. The resonant energy γ_r is given by $\gamma_r = (1 + a_w^2 + a_r^2 + P_z^2)^{1/2}$ with $P_z = \gamma_r \beta_z$ and the

synchronous velocity $\beta_z = k_r / (k_r + k_w)$. It follows that the relation between γ_r and γ_z is

$$\gamma_z = \gamma_r (1 + a_w^2)^{-1/2} \quad (4)$$

Hamiltonian (3) was obtained for small excursions $\tilde{\gamma}/\gamma_r \ll 1$ which is true for electrons not too far outside the separatrix. It is also assumed that the wiggler parameters change slowly compared to the wiggler wavelength $2\pi/k_w$. The term $\sin\psi_r$ parametrizes the rate of change for the resonant energy caused by the change in the wiggler wavelength

$$\frac{d}{dz}\gamma_r = - \frac{k_r a_w a_r}{\gamma_r} \sin\psi_r \quad (5)$$

where $\psi_r=0$ corresponds to an untapered wiggler. The terms proportional to $a_w a_r$ and $a_w a_s$ are the ponderomotive potentials due to the combined action of the wiggler with the main signal and the sideband respectively. The term δ_s in Eq. (4) signifies the Doppler shifted frequency departure of the sideband from the main signal

$$\delta_s = (k_w / k_r) (k_s - k_r). \quad (6)$$

When $a_s=0$ the Hamiltonian $H_0(P, \psi)$ is exactly integrable. The unperturbed trajectories, shown in Fig. 1(a), are given by

$$H_0(P, \psi) = K \quad (7)$$

where the constant K , regarded as the reduced energy in the ponderomotive frame, is determined from the initial conditions $K = H_0(P_0, \psi_0)$. These trajectories take the simplest possible form expressed in terms of the action angle variables (J, θ) defined as

$$J = \frac{1}{2\pi} \oint d\psi \tilde{\gamma}(K, \psi), \quad \theta = \frac{\partial}{\partial J} \int^\psi d\psi' \tilde{\gamma}(K, \psi'), \quad (8)$$

where $K = H_0(J)$ and the path of integration is over the unperturbed orbits. Clearly the action J is related to the area in phase space enclosed by the orbit. For untrapped particles and in case of untapered wiggler the limits of ψ integration are from 0 to 2π . In case of a tapered wiggler we may define the limits of integration starting at $\psi_s = 2\pi - \psi_r$ and going back to ψ_s around the separatrix. Thus J remains finite, avoiding an infinite jump in action across the separatrix that would otherwise result by considering the full orbit length for unbound orbits. In the above variables the width of the separatrix is given by J_s , the value of action at the separatrix.

We now apply the transformation (8) to the Hamiltonian Eq. (3) resulting into the decomposition of the perturbation $\cos[\psi(J, \theta) - \delta_s t]$ into harmonics of the synchrotron oscillation and obtaining

$$H(J, \theta) = H_0(J) + \sum_{n=0}^{\infty} \{ h_n^+(J) \cos(n\theta + \delta_s t) + h_n^-(J) \cos(n\theta - \delta_s t) \},$$

$$h_n^\pm(J) = \frac{a_w a_r}{\gamma_r} Q_n^\pm(J). \quad (9)$$

where $Q_n^\pm(J)$ are the Fourier coefficients for the synchrotron harmonics. This expansion is similar in spirit with the more familiar

case of expanding a plane wave propagating across a magnetized plasma into cyclotron harmonics. Expression (9) for the Hamiltonian is formally independent on the choice of tapered or untapered wiggler which will only affect the functional forms for $Q_n^\pm(J)$.

Here we present the analysis for constant parameter wiggler where the transformations are solvable in closed forms. The relation between new (J, θ) and old (P, ψ) variables is expressed by

$$\begin{aligned}
 J &= \begin{cases} J_s [E_2(\lambda) - (1-\lambda^2)E_1(\lambda)], & \lambda^2 < 1 \\ 2J_s \lambda E_2(1/\lambda), & \lambda^2 > 1 \end{cases} \\
 J_s &= \left(\frac{a_w a_r k_r}{2k_w} \right)^{1/2} \\
 \sin \frac{\psi}{2} &= \begin{cases} \lambda \operatorname{sn} \left[\frac{2}{\pi} E_1(\lambda) \theta \right], & \lambda^2 < 1 \\ \operatorname{sn} \left[\frac{1}{\pi} E_1(1/\lambda) \theta \right], & \lambda^2 > 1 \end{cases}
 \end{aligned} \tag{10}$$

with E_1, E_2 the complete elliptic integrals of the first and second kind and sn the Jacobi elliptic sine function. The trapping parameter λ^2 is given by

$$\lambda^2 = \frac{K + G}{2G}, \quad G = \frac{a_w a_r}{\gamma_r} \tag{11}$$

and we have $\lambda^2 < 1$, $\lambda^2 > 1$ for trapped and untrapped particles respectively. The three constants of the motion λ^2 , K and J are mutually dependent and any one of them defines uniquely a trajectory.

Since $H_0(J)$ is independent of θ the unperturbed trajectories in J, θ are straight lines

$$\begin{aligned}
J(z) &= \text{const.}, \\
\theta(z) &= \theta_0 + \kappa_b(J)z, \\
\kappa_b(J) &= dH_0(J)/dJ.
\end{aligned} \tag{12}$$

$\kappa_b(J)$ is the synchrotron wavenumber, connected to the synchrotron length L_b by $\kappa_b = 2\pi/L_b$ and to the synchrotron frequency $\omega_b(J)$ by

$$\kappa_b(J) = \frac{\omega_b(J)}{\beta_z c}. \tag{13}$$

From (8), (12) and (13) we find

$$\omega_b(J) = \begin{cases} \omega_b(0) \frac{\pi}{2E_1(\lambda)}, & \lambda^2 < 1 \\ \omega_b(0) \frac{\pi\lambda}{E_1(1/\lambda)}, & \lambda^2 > 1 \end{cases} \tag{14}$$

where

$$\omega_b(0) = \frac{c\beta_z}{\gamma_r} (a_w a_r k_w k_r)^{1/2}.$$

The coefficients $Q_n^\pm(J)$ are obtained by contour integration in the complex plane using the double periodicity properties of the Jacobi elliptic functions,

$$Q_n^\pm = -(\pm 1)^n \frac{n\pi^2}{E_1^2(\lambda)} \frac{q^{\frac{n}{2}}}{1 - (-q)^n}, \quad q = \exp\left(\frac{\pi E_1'(\lambda)}{E_1(\lambda)}\right), \quad \lambda^2 < 1 \tag{15}$$

$$Q_n^\pm = - \frac{n\pi^2 \lambda^2}{E_1^2(1/\lambda)} q^n \left(\frac{1}{1 - q^{2n}} \pm \frac{1}{1 + q^{2n}} \right), \quad q = \exp\left(\frac{\pi E_1'(1/\lambda)}{E_1(1/\lambda)}\right), \quad \lambda^2 > 1$$

$Q_n(J)$ tends to 0 as J tends to either 0 or ∞ while the maximum occurs at the separatrix $\lambda^2=1$. $Q_n(J)$ becomes progressively smaller with increasing n everywhere except in the vicinity of the separatrices. That limits the number of harmonics to be considered for instabilities localized away from the separatrix. The perturbed equations of motion in the standard form read

$$\begin{aligned} \frac{d}{dz} J &= \sum_{n=0}^{\infty} n \{ h_n^+(J) \sin(n\theta + \delta_s z) + h_n^-(J) \sin(n\theta - \delta_s z) \}, \\ \frac{d}{dz} \theta &= \omega_b(J) + \sum_{n=0}^{\infty} \left\{ \frac{dh_n^+}{dJ} \cos(n\theta + \delta_s z) + \frac{dh_n^-}{dJ} \cos(n\theta - \delta_s z) \right\}. \end{aligned} \quad (11)$$

The growth g_s for a given sideband is determined by the energy balance equation

$$\frac{d}{dz} a_s^2 = 2g_s a_s^2 = \frac{4\pi}{\omega_s} \{ \langle j_y(z, t) \frac{1}{2} (i a_s e^{ik_s z - i\omega_s t} + cc) \rangle + cc \} \quad (17)$$

The bracket $\langle \rangle$ denotes the average of a quantity $\langle f \rangle = L^{-1} \int f(z) dz$ over the short length scale $L = 2\pi/k_w \ll L_b = 2\pi/k_b$, and j_y is the transverse current

$$j_y(z, t) = en_b \int_0^{2\pi} d\theta \int_0^{\infty} dJ v_y \delta f(\theta, J; z) \quad (18)$$

with

$$v_y = \gamma^{-1} p_y = -(e/\gamma mc) A_y(z, t), \quad t(z) \approx z/c\beta_z.$$

For small sideband signal δf is the perturbed distribution f_1 under the Hamiltonian flow, Eqs. (16)

$$\frac{\partial f_1}{\partial z} + \kappa_b(J) \frac{\partial f_1}{\partial \theta} = - \frac{dJ}{dz} \frac{df_0}{dJ}, \quad (19)$$

where f_0 is the equilibrium distribution along the unperturbed orbits $df_0/dz=0$. Any distribution $f_0(J)$ depending on J alone is invariant under Eqs. (12)

$$\frac{d}{dz} f_0 = \frac{dJ}{dz} \frac{df_0}{dJ} = 0$$

thus f_0 is not uniquely defined.

We now assume a small imaginary part in the sideband wavenumber $g_s = \text{Im}(k_s)$. To obtain g_s for a given equilibrium $f_0(J)$ we solve Eq. (19) for f_1 , then substitute Eqs. (9) and (18) into (17) and finally average over the short length scales L . The computation is similar to that described in Ref. 13, with the exception that here the equations of motion and the evolution of the distribution function are parametrized by the length z rather than the time t , yielding

$$g^\pm = \mp k_r a_w^2 \frac{\omega_r}{\omega_s} \frac{\pi^2 \omega_p^2}{4 \gamma_r \omega_r^2} \sum_{n=1}^{\infty} \frac{|Q_n^\pm(J)|^2}{\gamma_r} \left(\frac{df_0}{dJ} \right)_{J_n} \left(\frac{d\omega_b}{dJ} \right)_{J_n}^{-1}, \quad (20)$$

with ω_p the beam plasma frequency. The superscript $+$ or $-$ corresponds to upper $\omega_s > \omega_r$ and lower $\omega_s < \omega_r$ sideband respectively. J_n is given implicitly by the resonant condition

$$\pm n \omega_b(J_n) = (k_w/k_r) (\omega_s - \omega_r), \quad (21)$$

as illustrated in Fig. 1(b). The sum over n on the right-hand side of Eq. (20) includes the contribution from all resonant groups of particles. The action J_n in Eq. (17) labels the orbit having the n th harmonic of the local bounce frequency in resonance with the sideband. The growth is determined by the slope of the distribution function f_0 near these resonant orbits $J \approx J_n$.

III. RESULTS AND CONCLUSIONS

The opposite signs in the right-hand side of Eq. (20) denote that upper and lower sidebands located symmetrically around the main signal $|\omega_s^+ - \omega_r| = |\omega_s^- - \omega_r|$ have opposite growth rates, $g^+/g^- = -1 + O(\delta\omega/\omega_r)^2$. The physical reason is connected with quantum mechanical considerations. The practical implication is that some sideband is always unstable except in case of trivial equilibrium $df_0/dJ=0$. It also reveals that, once destabilized, upper and lower modes have growth rates of equal magnitude and are equally dangerous. We clarify that, in general, the opposite signs do not imply that all upper sidebands have the same kind of stability, opposite to the stability of all lower sidebands. Depending upon f_0 stability may change sign between two upper (or lower) frequencies because the slopes df_0/dJ_n change as the location of the resonant J_n 's shift with ω_s .

For a monotonic distribution f_0 we observe that trapped and untrapped particles yield opposite contributions to a given mode because $d\omega_b/dJ$ changes sign across the separatrix. If trapped particles are stabilizing untrapped are destabilizing and vice versa.

High shear is stabilizing tending to reduce the magnitude of the growth g_s . This is expected as the number of resonant orbits is inversely proportional to $|d\omega_b/dJ|$. Shear tends to infinity near the separatrix thus the modes generated by electrons orbiting near the separatrix, corresponding to small $|\omega_s - \omega_r| \ll (k_r/k_w)n\omega_b(0)$, have the smallest growth rates. Modes coming from electrons near the center $J_n=0$ with $|\omega_s - \omega_r| \approx (k_r/k_w)n\omega_b(0)$ have also small growth rates as the coefficients $Q_n(J)$ tend to zero there. In particular, given any smooth distribution $f_0(J)$, electrons exactly at the bottom of the well have a null contribution to the instability. It takes a singular distribution of infinite gradient at $J=0$ to create an instability from electrons exactly at the center of the separatrix. The shear is generally higher for untrapped particles; thus the trapped ones usually dominate the instability.

Although Eq. (20) was obtained on purely classical arguments it nevertheless admits the correct quantum mechanical interpretation. Expressing $(df_0/dJ)(d\omega_b/dJ)^{-1}$ as $df_0/d\omega_b$ we observe that the growth g_s is proportional to the difference in population of oscillation quanta $\hbar\omega_b$ between the energy levels across the resonance. A more detailed explanation for the opposite signs in Eq. (20) is given in Ref. 13.

The normalized growth g_s/ω_r is plotted against the percentage mismatch $(\omega_s - \omega_r)/\omega_r$ for both upper and lower sidebands in Figs. 2 and 4. The contribution up to the third harmonic $n \leq 4$ in Eq. (20) is included in these plots. The parameters chosen correspond to a wiggler wavelength $\lambda_w = 3$ cm, $a_w = 5$, main signal strength $a_s = 5 \times 10^{-4}$, beam energy of 16.50 MeV ($\gamma = 32.3$) and current density $j = 100$ A/cm² (beam density 6.25×10^{10} cm⁻³). The equilibrium distribution is a Gaussian

$f_0(J) = (1/2\pi D^2)^{1/2} \exp(-(J-J_0)^2/2D^2)$ centered halfway inside the island, $J_0 = J_s/2$, and of width D equal to half the separatrix distance $D = J_s/2$. We plot the contribution of only the fundamental, $n=1$ in Eq. (20), in Fig. 2(a) adding the first harmonic $n=2$ in (b) the second harmonic in (c) and the third harmonic in (d). New unstable bands emerge with each harmonic, while the growth for already unstable bands is modified. For example we observe two upper and two lower unstable bands in (b) but only one upper and three lower bands in (d). We find the contributions from higher than the third harmonics $n > 4$ generally negligible. The upper frequency $\omega_s > \omega_r$ and the lower frequency $\omega_s < \omega_r$ parts of the unstable spectrum come from the regimes of negative and positive slope $df_0/d\omega_b$ respectively, shown in Fig. 3. The lower sideband growth is peaking at frequencies corresponding to $n\omega_b(J_{\max})$ with $J_{\max}$ the value maximizing $df_0/d\omega_b(J)$. The peaks for the upper sideband growth, however, do not occur at $J=0$ that minimizes $df_0/d\omega_b(J)$, but at J halfway inside the negative slope regime. This is because $Q_n(J)$ and consequently g_s are zero at $J=0$, showing the negligible contribution from electrons at the bottom of the well. The analogous effect in plasma physics is the elimination of the thermal effects when ρ_L goes to zero. In any case the most unstable modes are far from the frequencies $|\omega_s - \omega_r| = (k_r/k_w)n\omega_b(0)$ pointed by the arrows in Fig. 2(d).

Since sidebands can not be shut out completely it remains debatable whether a distribution function can be tailored experimentally to minimize their growth rate. From the previous discussion a flat distribution inside the trapped regime with sharp gradients localized at the separatrix seems the appropriate choice. Instabilities will then localize near the separatrix and the growth

will be suppressed by the strong shear. To check this we plot the growth rates from Eq. (20) for two types of distributions $f_0(J)$: (i) Two Gaussians $f_0(J) = (1/2\pi D^2)^{1/2} \exp(-J^2/2D^2)$ centered at the centre of the island and of characteristic lengths D equal to half the island width $D=J_s/2$ in Fig. 4(a) and one island width $D=J_s$ in Fig. 4(b). (ii) Two "step-like" distributions of the form $f_0(J) = (1/\alpha D)^{1/2} \exp(-(J/\alpha D)^N)$ with $N=16$. Selecting $\alpha = (N/N-1)^{1/N}$ places the sharp gradient at $J=D$ and we plot the case $D=J_s/2$ in Fig. 4(c) and $D=J_s$ in Fig. 4(d).

Comparing Fig. 4(a) to Fig. 4(b) and Fig. 4(c) to Fig. 4(d) it is seen that the growth rates between similar types of distributions tend to decrease as the location of the maximum gradient $(df_0/dJ)_{\max}$ approaches the separatrix. In both cases there is more than one order of magnitude reduction in the growth by shifting the maximum gradient position from $D=J_s/2$ to $D=J_s$. Because f_0 was chosen monotonic in all above plots and because it was limited to trapped particles, $df_0/d\omega_b$ preserves sign and only lower modes are unstable. The spectral width of the unstable regimes is reduced with a parallel increase in the maximum growth as one goes from the Gaussian type to the step-like type of distributions. Also the distance of the sideband frequencies from the main signal decreases by shifting the gradient position D closer to the separatrix. Distributions with sharp gradients at the separatrix such as those in Figs. 4(b) and 4(d) are perhaps more relevant to the case of variable wiggler FEL, where the "bucket" of the trapped particles is decelerating in phase space leaving the untrapped particles behind, and only small diffusion occurs across the separatrix allowing sharp gradients.

The limit of a δ -function distribution $f_0 = \delta(J-J_0)$, examined

elsewhere¹³, is the most unstable case but of the least practical interest, because even the case of a monoenergetic beam distribution $p_z = p_0$ is described in J-space by a smooth distribution $f_0(J)$ of finite width ΔJ (see Fig. 1).

Our calculations done for the case of a fixed wavelength wiggler can be easily extended to the case of a variable wiggler FEL provided that the same adiabatic assumptions hold. Only the functional relations Eqs. (7) between the action-angle variables (J, θ) and the coordinates (P, ψ) need to change. The derivation of the growth rate, performed in J-space, is independent of the transformation $J(P, \psi)$ and the result, Eq. (20), stands as is. It is the high growth of tapered wiggler FELs that seems to challenge the adiabatic approach to the problem. In this case the change of the signal amplitude $a_r(t)$ in time together with the dependence of the equilibrium f_0 on both time and θ should be included for a more realistic treatment.

Acknowledgements

This work was supported by the U. S. Army Strategic Defense Command.

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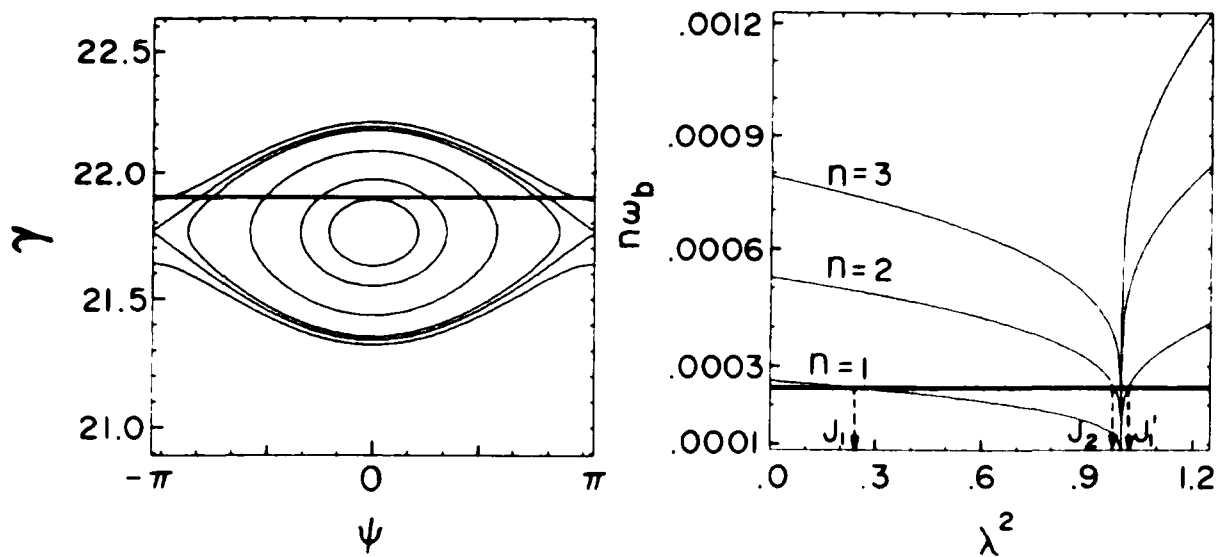


Figure 1. Time averaged motion without the sidebands. (a) Plots in phase space of the unperturbed orbits $H_0(P, \psi) = K$. The intersections with the horizontal line $P = \text{const.}$ mark the initial conditions for each orbit. (b) The normalized bounce frequency ω_b and the first two harmonics as functions of the trapping parameter $\lambda^2(J)$. The intersections with the horizontal line $\omega = (\omega_s - \omega_r)k_w/k_r$ determine the position J_n of the resonant orbits for a given ω_s .

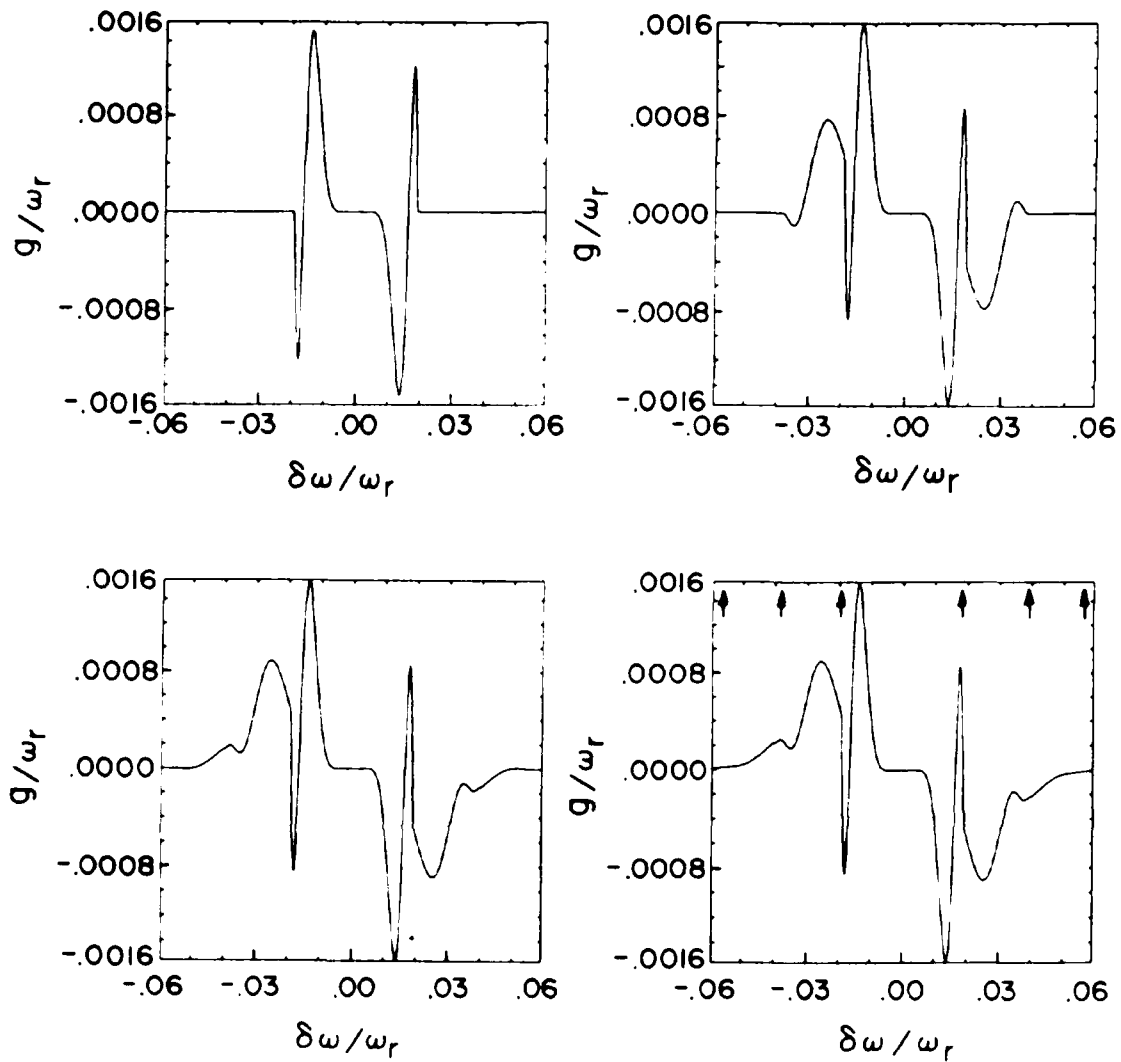


Figure 2. Growth for a Gaussian distribution $f_0(J) = C \exp(-(J - J_0)^2 / 2D^2)$ centered halfway inside the trapped particle island $J_0 = J_s/2$ and of width D equal to J_s . The normalized growth g_s/ω_r is plotted versus $\delta\omega/\omega_r$ for (a) the fundamental contribution $n=1$ in Eq. (20), (b) including the first harmonic $n=2$, (c) two harmonics and (d) three harmonics. The most unstable modes do not correspond to harmonics of the bounce frequency at the bottom $\omega_b(0)$ indicated by the arrows.

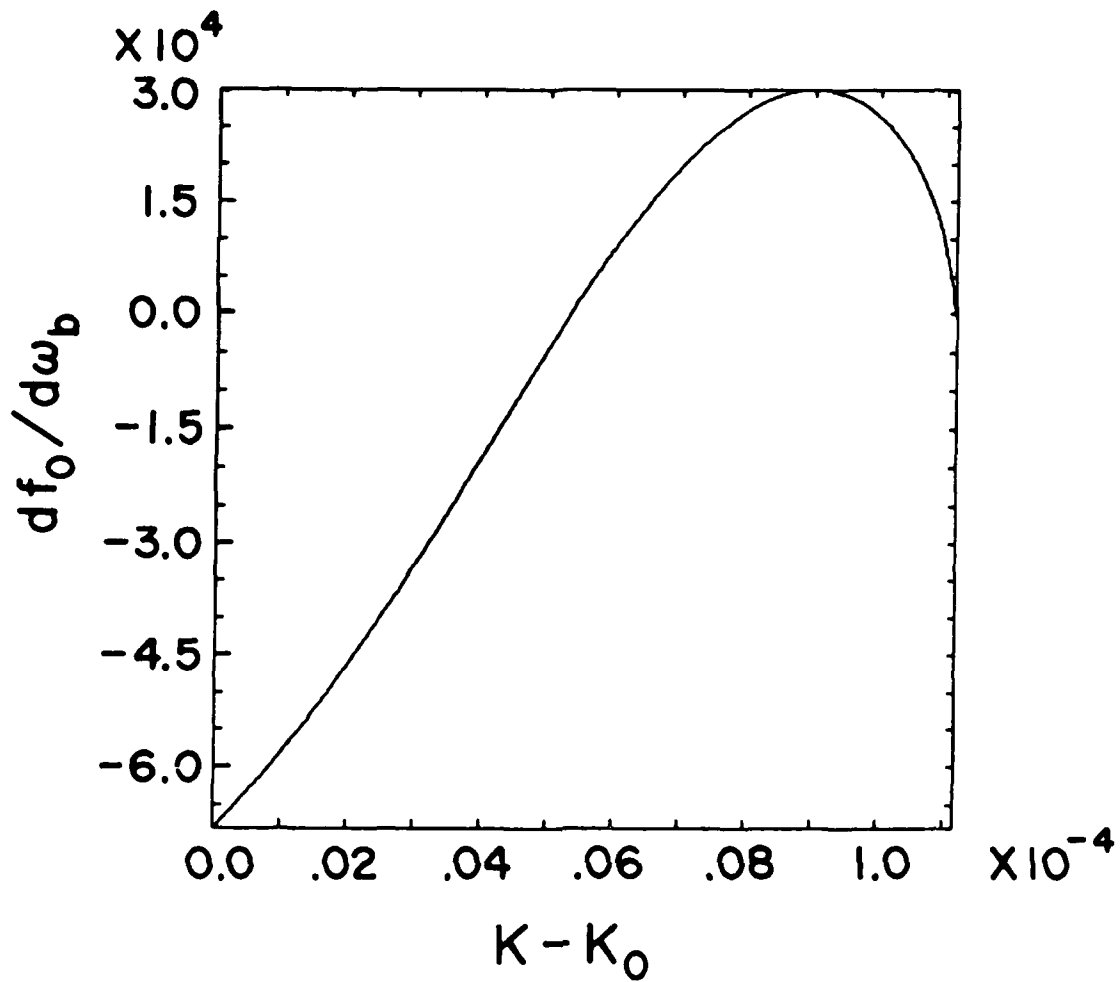


Figure 3. Plot of $df_0(J)/d\omega_b(J) = df_0/dJ(d\omega_b/dJ)^{-1}$ as a function of $K(J) - K(0)$ for the distribution f_0 of Figure 2. The slope goes to zero near the separatrix $K(J_s)$ because of the infinite shear $d\omega_b/dJ$.

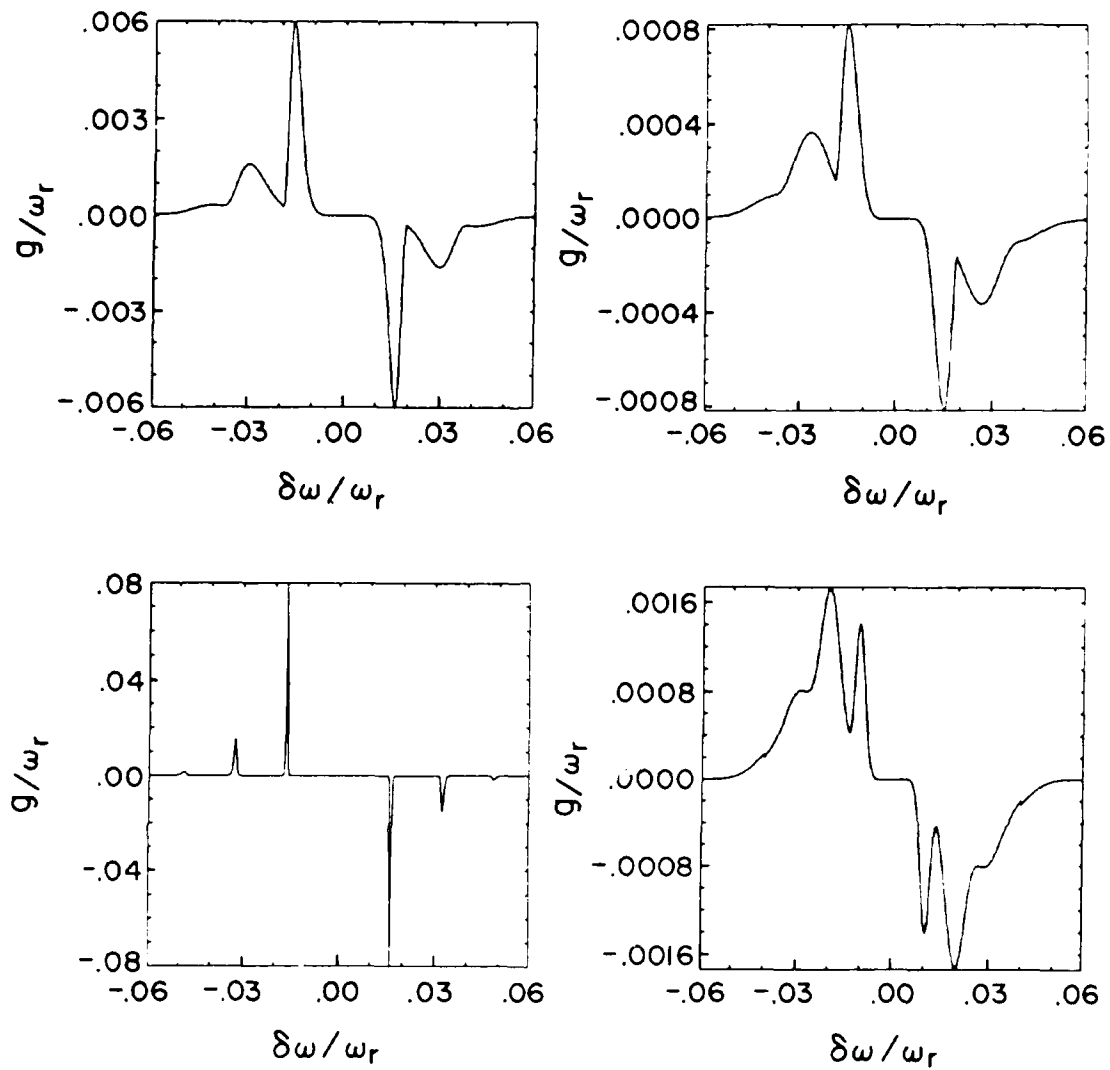


Figure 4. Normalized growth for monotonic distributions centered at the bottom of the well $J=0$ including the first three harmonics $n \leq 4$ in Eq. (20). (a) Gaussian distribution of width D equal to half the island width $D=J_s/2$ (b) Gaussian distribution with $D=J_s$. (c) Step-like distribution with $D=J_s/2$ and (d) Step-like distribution with $D=J_s$.

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