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Synchrotron-Betatron Parametric Instability in Free-Electron Lasers

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<p>Analysis of the synchrotron-betatron coupling in a free-electron laser indicates that the particle motion is susceptible to an instability which is of a fundamentally different physical character to those examined previously. The domain of this instability, which is of a parametric type and causes an exponential growth of the betatron amplitude, is identified and the growth rate therein determined.</p>			
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SYNCHROTRON-BETATRON PARAMETRIC INSTABILITY IN FREE-ELECTRON LASERS

I. Introduction

All electron beams have finite emittance; that is, there is a spread in the particle velocities transverse to the primary direction of motion. On traversing a wiggler, the inevitable presence of transverse gradients in the magnetic field lead to betatron oscillations of the particles. On the other hand, the beating of the wiggler and optical fields forms the ponderomotive potential wherein the particles perform synchrotron oscillations. Coupling of these two basic oscillatory degrees of freedom in a free-electron laser (FEL) may be deleterious to its efficient operation due either to detrapping out of the ponderomotive potential well or growth of betatron oscillation amplitude and emittance.

In Ref. 1 it is shown that the curvature of the optical wavefronts couples the synchrotron and betatron oscillations of the electrons and that under the proper resonance condition the amplitude of the synchrotron motion increases unstably. In subsequent work, the authors of Ref. 2 showed that a similar betatron-induced forcing of the synchrotron motion results in the case of planar optical wavefronts provided the wiggler field is tapered.

In what follows it is shown that the particle motion is additionally susceptible to an instability which is dependent neither on the curvature of the wavefronts nor on a tapered wiggler field. The instability arises from the energy-dependence of the betatron wavenumber. As a result when the electrons undergo synchrotron oscillations the betatron wavenumber is modulated at the synchrotron period, resulting in a parametric instability. This leads to an exponential growth of the betatron oscillation amplitude

and velocity. The region of parameter space wherein this instability is operative is identified and an analytical expression for the growth rate is obtained.

II. Wiggle-average equations of motion

It is assumed that the wiggler is linearly polarized, with the vector potential given by

$$A_w = A_w \cosh(k_w y) \sin(k_w z) e_x,$$

where A_w is constant, $2\pi/k_w$ is the wiggler period and e_x is the unit vector along the x-axis. The vector potential of the linearly-polarized optical field is taken to correspond to the fundamental vacuum Gaussian mode:

$$A_s = A_s \cos\left(\frac{\omega}{c} z - \omega t + \phi\right) e_x,$$

where ω is the angular frequency, c is the speed of light in vacuo,

$$\phi(y, z) = \omega y^2 / (2cR) - \frac{1}{2} \arctan(z/z_R) + \phi_0, \quad (1a)$$

$$R = z \left(1 + (z_R/z)^2 \right) \quad (1b)$$

is the radius of curvature of the wavefronts,

$$z_R = \frac{\omega w_0^2}{2c} \quad (1c)$$

is the Rayleigh range, ϕ_0 is a constant,

$$A_s = \hat{A}_s \left(\frac{w_0}{w(z)} \right)^{1/2} \exp\left(-y^2/w^2(z)\right), \quad (1d)$$

\hat{A}_s is a constant, and $w(z)$, the spot size at z , is related to the waist w_0 at $z = 0$ by

$$w(z) = w_0 \left[1 + (z/z_R)^2 \right]^{1/2}. \quad (1e)$$

It should be remarked that the amplitude and phase of the optical field are assumed to be constant in time.

Following the standard procedure,³ the equations of motion in terms of the independent variable z are derivable from the Hamiltonian function $-p_z(y, p_y; t, -E; z)$, which is given by

$$p_z = \frac{E}{c} - \frac{m^2 c^3}{2E} \left(1 + (p_y/mc)^2 + a_x^2 \right), \quad (2)$$

where E , the total energy of an electron of rest-mass m and charge $-|e|$, is canonically conjugate to the time t , p_y is the momentum conjugate to the coordinate y , and a_x is the normalized, total vector potential:

$$a_x = a_w \cosh(k_w y) \sin(k_w z) + a_s \cos\left(\frac{\omega}{c} z - \omega t + \phi\right), \quad (3)$$

with $a_{w,s} = |e| A_{w,s} / (mc^2)$. Since the vector potential is not an explicit function of the coordinate x , the associated canonical momentum P_x is a constant. Equation (2) has been derived by choosing $P_x = 0$, and omitting terms on the order of $(mc^2/E)^{-2}$ and higher in the expansion of the Hamiltonian.

Upon squaring the expression in Eq. (3), averaging over the wiggler period and retaining only the slowly-varying contribution to the ponderomotive potential, it is found that

$$p_z = \frac{E}{c} - \frac{m^2 c^3}{2E} \left\{ 1 + (p_y/mc)^2 + \frac{1}{2} a_w^2 \cosh^2(k_w y) + a_w a_s \cosh(k_w y) \right. \\ \left. \times \sin\left[\left(\frac{\omega}{c} + k_w\right)z - \omega t + \phi\right] \right\}.$$

Hamilton's equations are then given by

$$\frac{dy}{dz} = v, \quad (4a)$$

$$\frac{dv}{dz} = -k_{\beta}^2 \left[\frac{\sinh(2k_w y)}{2k_w} \right], \quad (4b)$$

$$\frac{d\psi}{dz} = k_w - \frac{\omega}{2\gamma^2 c} \left[1 + \frac{1}{2} a_w^2 \cosh^2(k_w y) + (\gamma v)^2 \right], \quad (4c)$$

$$\frac{d\gamma}{dz} = \frac{-\omega}{2\gamma c} a_w a_s f_B \cosh(k_w y) \cos(\psi + \phi), \quad (4d)$$

where $\gamma = E/(mc^2)$, $\psi = (\omega/c + k_w)z - \omega t$, $v = p_y/(\gamma mc)$, $f_B = J_0(\xi) - J_1(\xi)$ is the usual difference of Bessel functions,⁴ $\xi = (a_w/2)^2/(1 + a_w^2/2)$, and

$$k_{\beta} = \frac{a_w k_w}{\sqrt{2} \gamma} \quad (5)$$

is the betatron wavenumber. In arriving at Eq. (4b) it has been assumed that

$$\max \left\{ \frac{a_s}{a_w}, \frac{\gamma a_s}{1 + a_w^2/2}, \frac{a_s}{a_w (k_w w_0)^2} \right\} \ll 1, \quad (6)$$

where w_0 is the waist of the optical field, Eq. (1e). Additionally, the wave term, proportional to a_s/γ^2 , has been dropped from Eq. (4c).

III. Analysis of synchrotron-betatron coupling

Variation of the optical phase ϕ across the wavefronts underlies the instability investigated in Ref. 1. Specifically, the first term in (1a) leads to a linear increase with z of the synchrotron oscillation amplitude of deeply trapped electrons. Assuming the radius of curvature R of the wavefronts to be constant, the growth rate of this instability is proportional to y_{β}^2/R , where y_{β} is the amplitude of the betatron

oscillation. In what follows it is assumed that $y_{\beta}^2/R \rightarrow 0$. This is the case for electrons confined to the central portion of the electron beam, or in the region where the curvature of the wavefronts is small. Assuming further that the z-variation of ϕ is small on the scale length of the instability to be discussed, for a synchronous electron Eqs. (4c) and (4d) imply

$$\gamma = \hat{\gamma} = \text{constant}, \quad (7a)$$

$$\hat{\psi} + \phi = (n + \frac{1}{2})\pi, \quad n = 0, \pm 1, \pm 2, \dots \quad (7b)$$

whence, for $|k_w y| \ll 1$, Eqs. (4a) and (4b) integrate to

$$y = y_{\beta} \cos(\hat{k}_{\beta} z) + \hat{k}_{\beta}^{-1} v_{\beta} \sin(\hat{k}_{\beta} z), \quad (8)$$

where, from (5), $\hat{k}_{\beta} = a_w k_w / \sqrt{2} \hat{\gamma}$ is the betatron wavenumber for a synchronous electron, and y_{β} , v_{β} are constants. Upon substituting (8) into Eq. (4c) one obtains

$$\begin{aligned} \frac{1}{2} a_w^2 \cosh^2(k_w y) + (\gamma v)^2 &= \frac{1}{2} a_w^2 [1 + (k_w y)^2] + (\gamma v)^2 \\ &\rightarrow \frac{1}{2} a_w^2 + \hat{\gamma}^2 (k_{\beta}^2 y_{\beta}^2 + v_{\beta}^2), \end{aligned}$$

which is a constant in z, whence the synchronous energy follows from Eq. (4c) by setting $d\hat{\psi}/dz = 0$:

$$\hat{\gamma}^2 = \frac{1 + \frac{1}{2} a_w^2 [1 + (k_w y_{\beta})^2]}{2ck_w/\omega - v_{\beta}^2}. \quad (9)$$

For an electron in the vicinity of the synchronous particle, $\gamma = \hat{\gamma} + \delta\gamma$, $\psi = \hat{\psi} + \delta\psi$, and Eqs. (4c) and (4d) imply synchrotron oscillations of $\delta\gamma$ and of $\delta\psi$. Through the γ -dependence of the betatron

wavenumber, Eq. (5), it follows from Eqs. (4a) and (4b) that under an appropriate resonance condition a parametric instability of the coupled betatron and synchrotron oscillations is then possible.

To examine this, the expression in (8) is generalized to

$$y = y_{\beta}(z)\cos[(\hat{k}_{\beta} + \epsilon)z] + (\hat{k}_{\beta} + \epsilon)^{-1}v_{\beta}(z)\sin[(\hat{k}_{\beta} + \epsilon)z], \quad (10)$$

where y_{β} and v_{β} are now functions of z and ϵ is a small wavenumber shift to be determined. Assuming $|dy_{\beta}/dz| \sim \epsilon y_{\beta}$, $|dv_{\beta}/dz| \sim \epsilon v_{\beta}$, and omitting terms on the order of ϵ^2 , substitution of (10) into Eqs. (4a) and (4b) yields

$$\begin{aligned} & (dv_{\beta}/dz - \epsilon \hat{k}_{\beta} y_{\beta})\cos[(\hat{k}_{\beta} + \epsilon)z] - (\hat{k}_{\beta} dy_{\beta}/dz + \epsilon v_{\beta})\sin[(\hat{k}_{\beta} + \epsilon)z] \\ & = \hat{k}_{\beta}^2 \left\{ y_{\beta}\cos[(\hat{k}_{\beta} + \epsilon)z] + (\hat{k}_{\beta} + \epsilon)^{-1}v_{\beta}\sin[(\hat{k}_{\beta} + \epsilon)z] \right\} (\delta\gamma/\hat{\gamma}). \end{aligned} \quad (11)$$

To obtain an expression for $\delta\gamma(z)$ to be used in Eq. (11), Eqs. (4c) and (4d) are perturbed about the synchronous values (7a) and (7b):

$$\begin{aligned} \frac{d}{dz} \delta\psi &= k_w - \frac{\omega}{2c\hat{\gamma}} \left[1 + \frac{1}{2} a_w^2 (1 + k_w^2 y^2) + \hat{\gamma}^2 v^2 \right] \\ &+ \frac{\omega}{c\hat{\gamma}^3} \left[1 + \frac{1}{2} a_w^2 (1 + k_w^2 y^2) \right] \delta\gamma, \end{aligned} \quad (12a)$$

$$\frac{d}{dz} \delta\gamma = (-1)^n \frac{\omega a_w a_s f_B}{2\hat{\gamma}c} \left(1 + \frac{1}{2} k_w^2 y^2 \right) \delta\psi, \quad (12b)$$

where n denotes an integer from the set defined in (7b). Substituting (10) and the forms

$$\delta\gamma = \Delta\gamma(z)\cos[2(\hat{k}_{\beta} + \epsilon)z + \eta], \quad (13a)$$

$$\delta\psi = \Delta\psi(z)\sin[2(\hat{k}_{\beta} + \epsilon)z + \eta], \quad (13b)$$

where η is a constant, into (12) and assuming that

$$k_w y_\beta \ll \hat{\gamma}^{-1/2}, \quad (14)$$

one obtains an identity whose consistency requires that $\Delta\gamma$ and $\Delta\psi$ be independent of z and that

$$2(\hat{k}_\beta + \epsilon) = \hat{k}_{\text{syn}}, \quad (15)$$

where

$$\hat{k}_{\text{syn}} = \left[\frac{\omega^2 a_w a_s f_B}{2\hat{\gamma}^4 c^2} \left(1 + \frac{1}{2} a_w^2 \right) \right]^{1/2}$$

is the synchrotron wavenumber. In addition, it is required that the integer n in Eq. (12b) be odd. It should be noted that the constraint indicated by (14) is consistent with the neglect of the radiation field term in Eq. (4c) and the parameter restriction in (6).

Now substituting (13a) into Eq. (11) one again obtains an identity. Assuming $y_\beta, v_\beta \sim \exp(\lambda z)$, the consistency condition reduces to

$$\lambda = \pm \left\{ \left[\frac{1}{2} \hat{k}_\beta (\Delta\gamma/\hat{\gamma}) \cosh\eta \right]^2 - \epsilon^2 \right\}^{1/2}, \quad (16)$$

which is the growth rate of the instability.

IV. Discussion

It is clear that the source of the instability discussed here lies in the energy dependence of the betatron wavenumber, Eq. (5). Previous studies of the synchrotron-betatron coupling have neglected this dependence. The results of Refs. 1 and 2 indicate a linear growth in the amplitude of the synchrotron motion ($\Delta\gamma, \Delta\psi$) for deeply trapped electrons.

On the other hand, for the parametric instability discussed herein, $\Delta\gamma$ and $\Delta\psi$ are constants, whereas y_β and v_β grow exponentially fast. In all three cases, however, the resonance condition is essentially $\hat{k}_{\text{syn}} \approx 2\hat{k}_\beta$.

From Eqs. (15) and (16) it follows that for a sufficiently large wavenumber mismatch ϵ the growth rate vanishes. Maximum growth is obtained for $\epsilon = 0$:

$$\lambda_{\text{max}} = \frac{1}{2} \hat{k}_\beta |(\Delta\gamma/\hat{\gamma}) \cosh|. \quad (17)$$

Since, at the order of approximation employed here, $\Delta\gamma$ and $\Delta\psi$ are constants, the parametric instability is expected to affect the extraction efficiency of the FEL mainly through increased electron beam radius and particle detrapping from the ponderomotive potential, such as that described in Refs. 1 and 2, would result once the betatron amplitude is sufficiently large. It must be remarked, however, that if the variation of the optical field amplitude or phase over the scale length indicated by (17) is large enough, then the resonance (15) may be crossed without an appreciable effect on the efficiency.^{5,6}

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