# Generation of ultrashort radiation pulses by injection locking a regenerative free-electron-laser amplifier

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We demonstrate how a steady-state train of ultrashort radiation pulses can be produced utilizing a new free-electron laser (FEL) configuration, the injection-locked regenerative klystron amplifier (IRKA). This configuration consists of two elements: (1) a prebuncher, which microbunches the electron beams at the desired output wavelength, and (2) a multipass FEL operated at a very small cavity desynchronism and below the lasing threshold, in the regime of regenerative amplification. The regenerative amplifier is driven by the microbunched electron beam, so that the pulse-to-pulse stability is provided by the pre-buncher. The broad amplification bandwidth of this regenerative amplifier enables generation of ultrashort pulses, much shorter than a slippage length, with high efficiency. The IRKA configuration can produce such ultra-short radiation pulses while avoiding the chaotic dynamics that limits conventional FEL performance. [S1063-651X(97)12609-5]

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## I. INTRODUCTION AND MOTIVATION

The free-electron laser (FEL) can produce tunable high-power radiation. Another potentially important feature of the FEL is its ability to generate ultrashort radiation pulses [1]. As FEL technology matures, it becomes more important to provide the radiation users with some of the capabilities that more conventional lasers have acquired over decades of research and development. For example, many scientific applications, such as condensed-matter studies and chemical dynamics, require ultrashort pulses (of order 100 fs in duration), with precise frequency and timing control, and rapid tunability. In solid-state laser systems, such capability is typically achieved by using the output of a low-power frequency-stable oscillator to injection-lock a high-power, broadband regenerative amplifier. As this paper demonstrates, a similar approach can be used with free-electron lasers. Injection-locking of a regenerative FEL amplifier may become a useful tool for generating short, high-power, tunable radiation pulses, with low pulse-to-pulse jitter. Short duration radiation pulses are important for studying ultrafast phenomena. Small pulse-to-pulse jitter, or steady-state operation, is essential for high repetition rate pump-probe experiments. While separately each of these two characteristics, ultrashort pulse (with radiation pulses much shorter than the slippage length) and stable operation, has already been experimentally demonstrated, the challenge still remains to develop a single FEL device which has both features.

Most FEL user facilities [2–4] operate as low gain oscillators and the majority are driven by radio-frequency (rf) linacs. The electron beam consists of very short (picosecond or subpicosecond) electron bunches, spaced by a multiple of the rf wavelength. The radiation pulse length is determined by the bunch length $L_b$, if $\Delta \ll L_b$, where $\Delta = N_w \lambda$ is the slippage length (the distance the optical pulse outruns the electron bunch in one pass), $\lambda$ is the radiation wavelength, and $N_w$ is the number of wiggler periods. For electron bunches much shorter than the slippage distance, the pulse length is (roughly) determined by the slippage length, $\Delta$. To make more precise statements about the duration and structure of radiation pulses in short-bunch oscillators requires the inclusion of cavity desynchronism and losses. Cavity desynchronism is introduced to counter the tendency of the amplified radiation to fall behind the electron bunch after a number of passes (laser lethargy). Making the cavity slightly shorter than required for perfect synchronism with electron bunches in the absence of the FEL interaction restores the longitudinal overlap between radiation and electrons.

The supermode theory of low-gain FEL oscillators in the short-bunch limit was recently developed [5]. Many of the analytical conclusions parallel those of the pioneering numerical work of Dattoli *et al.* [6] on the linear radiation supermodes in oscillators driven by electron bunches with an arbitrary temporal profile. The equations describing the temporal pulse evolution of a short-bunch oscillator are very similar to those describing backward wave oscillators (BWO's) (see, for example, Ref. [7]). This implies that boundary conditions are specified at the head and the tail of the radiation pulse and is essential to our analysis of the pulse-shaping concept described in this paper.

Here we propose a scheme that circumvents an important limitation of short-pulse FEL oscillators. Very short radiation pulses can only be obtained in the operating regime where cavity desynchronism is very small. Radiation pulses as short as six wave periods (and about one sixth of the slippage length) were recently obtained experimentally [1]. The experimental measurement was done in the *linear* regime. Unfortunately, it is precisely in this regime of small cavity desynchronism that *nonlinear* multimode effects, such as limit-cycle saturation and transition to chaos, dominate the FEL dynamics. They result in nonrepetitiveness of radiation pulses from pass to pass after the oscillator enters the non-

 

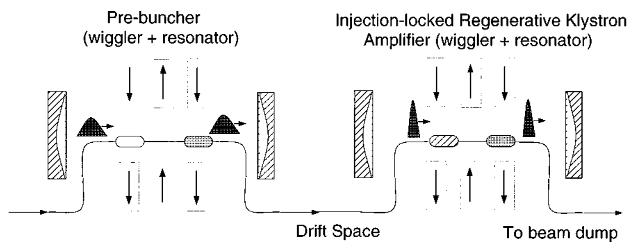


FIG. 1. Schematic of the injection-locked regenerative klystron amplifier (IRKA) with a FEL oscillator prebuncher. The shading of the electron beam indicates the degree of bunching. Radiation pulses in IRKA are much shorter than in the prebuncher.

linear regime. The proposed scheme enables the steady-state generation of ultrashort radiation pulses with high efficiency (i.e., larger than the typical value of $1/2N_w$).

The key idea is to operate a FEL as an injection-locked regenerative klystron amplifier (IRKA). The IRKA consists of a magnetostatic wiggler placed inside a high-$Q$ radiation cavity with a very small desynchronism. This cavity is operated *below* the lasing threshold, i.e., the FEL is run as a multipass amplifier. The seed current at the desired wavelength is provided by a prebuncher, which can be designed in a variety of ways. For example, masked chicanes can be utilized to prebunch the electrons [9]. Alternatively, a free-running FEL oscillator, operated above the lasing threshold and nonlinearly saturated in steady state, can be used as shown in Fig. 1. Unlike an IRKA, the prebunching oscillator would then operate with a relatively large cavity desynchronism, ensuring that there is only one *linearly unstable* radiation supermode which saturates quickly and at low intensity. An important difference between the IRKA and a traditional regenerative amplifier is that the latter is injection locked by radiation while the former is locked by a microbunched beam (thus making it a *klystron* amplifier).

The IRKA configuration is analyzed using the formalism of Ref. [5], which is briefly reviewed in Sec II. Section III describes the IRKA scheme and provides a numerical example of its implementation.

## II. THEORETICAL BACKGROUND

Our theoretical model is one-dimensional, low gain, and we assume an electron bunch length that is much less than a slippage length. The evolution of laser pulses is then controlled [5] by three parameters: the reduced Colson parameter, $\Gamma = j_C L_b/\Delta$, the cavity desynchronism $\delta\mathcal{L}$, and the cavity losses $\alpha_0$, where

$$j_C = 4\pi \frac{I_b f}{I_A}\left(\frac{N_w}{\gamma}\right)^3 \left(\frac{a_w \lambda_w F}{r_b}\right)^2. \tag{1}$$

Here $I_b$ is the electron current, $I_A = mc^3/e = 17$ kA is the Alfven limiting current, $f$ is the filling factor describing the transverse overlap between the optical and electron pulses, $\gamma$ is the relativistic factor of the beam electrons, $r_b$ is the beam radius, $\lambda_w = 2\pi/k_w$ is the wiggler period, $a_w = eB_w/k_w mc^2$ is the normalized vector potential of the wiggler, and $F$ is 1 for a helical wiggler and equal to the well-known difference of Bessel functions for a linear undulator. Moreover, by an appropriate normalization [5] of "time" (a continuous equivalent of the pass number), cavity desynchronism, and cavity losses

$$\tau = \Gamma n, \quad \nu = \frac{2\delta\mathcal{L}}{\Delta\Gamma}, \quad \alpha = \frac{\alpha_0}{\Gamma},$$

the evolution of a FEL oscillator in the limit of $L_b \ll \Delta$ is governed only by $\nu$ and $\alpha$.

The independent variables are $\tau$ and $\xi = (ct - z)/\Delta$, the normalized distance from the head of the radiation pulse (which is assumed to move with the speed of light $c$). The field is described by a slowly varying complex amplitude, $A$, normalized so that $|A|^2 = 4\pi N_w j_c P(z,t)/P_e$, where $P(z,t)$ is the intracavity optical power and $P_e = mc^2\gamma_0(I_b/e)$ is the electron beam power. Electrons are characterized by their ponderomotive phase $\theta = (k + k_w)z - \omega t$ and normalized momentum $p = d\theta/dz$, where $\omega/c = k = 2\pi/\lambda$.

As was first demonstrated in Ref. [8], the essential FEL physics can be captured by following the evolution of four complex moments of the electron distribution function: $B_1 = \langle e^{-i\theta}\rangle$, $P_1 = \langle p e^{-i\theta}\rangle$, $Q = \langle p\rangle$, and $\sigma^2 = \langle (p - Q)^2\rangle$, which have the physical meaning of bunching, momentum bunching, average momentum (also called momentum detuning), and momentum spread, respectively. These moments are functions of $\zeta$, which plays the role of a distance along the wiggler. Since fresh electrons enter the oscillator each pass, the moments of the distribution function depend on normalized pass number $\tau$ parametrically through the optical field. This truncated description of the electrons by four global moments [5,8] accurately describes such complex FEL phenomena as nonlinear saturation of FEL amplifiers, synchrotron oscillations, nonlinear superradiance, and limit cycles in FEL oscillators.

The reduced dynamics of a FEL oscillator driven by ultrashort electron bunches is described by a set of equations [5]

$$\frac{\partial A}{\partial \tau} - \nu \frac{\partial A}{\partial \xi} + \frac{\alpha}{2} A = \eta B_1, \tag{2}$$

$$\frac{\partial B_1}{\partial \xi} = -iP_1, \tag{3}$$

$$\frac{\partial P_1}{\partial \xi} = -A - iSB_1 - 2iQP_1 + 2iQ^2 B_1, \tag{4}$$

$$\frac{\partial Q}{\partial \xi} = -[AB_1^* + \text{c.c.}], \tag{5}$$

$$\frac{\partial S}{\partial \xi} = -2[AP_1^* + \text{c.c.}], \tag{6}$$

where $S = \sigma^2 + Q^2 = \langle p^2\rangle$, $\eta(\xi) = 1$ if $0 < \xi < 1$ and $\eta(\xi) = 0$ otherwise. Equations (3)–(6) are integrated between $\xi = 0$ and 1. Physically, $\xi = 0$ corresponds to the radiation slice which overlaps the electron bunch at the entrance into the wiggler while $\xi = 1$ corresponds to the radiation slice which overlaps the electron bunch at the exit of the wiggler. Radiation which has already slipped ahead of the electron bunch due to finite cavity desynchronism $\nu>0$ occupies the region $\xi<0$, where it decays exponentially as $\exp(\alpha\xi/2\nu)$.

We consider $\nu>0$ (shortened cavity), for which causality requires that the boundary condition for the radiation amplitude be $A(\xi=1,\tau)=0$. Boundary conditions for particle variables are fixed at the wiggler entrance, i.e., at $\xi=0$. Setting boundary conditions at different boundaries for particles and fields is reminiscent of equations that describe backward wave oscillators [7], where particles are initialized at the *entrance* into the interaction region while the microwave fields (which, due to the slow-wave structures, have a group velocity in the opposite direction to their phase velocity) are initialized at the *exit* of the interaction region. Initializing particles and fields at different locations is intuitively clear for a backward wave oscillator, but is less intuitive for FEL oscillators, where the ''interaction window'' $0<\xi<1$ itself moves with the speed of light, and is equal to one slippage length.

Equations (2)–(6) can be linearized by assuming that $Q$ and $S$ are constant [that is, by discarding Eqs. (5)–(6)]. A further simplifying assumption is that electron bunches enter the wiggler with the same level of prebunching, momentum detuning, and momentum spread from pass to pass. Thus

$$B_1(\xi=0,\tau)=B_{10}, \quad P_1(\xi=0,\tau)=P_{10},$$
$$Q(\xi=0,\tau)=y_0, \quad S(\xi=0,\tau)=y_0^2+\sigma_0^2. \tag{7}$$

The solutions of the linearized equations can be written as a sum of a stationary solution with *inhomogeneous* boundary conditions, given by Eq. (7) at $\xi=0$, and a series of time-dependent solutions with *homogeneous* (i.e. vanishing) boundary conditions. For example,

$$A(\xi,\tau)=A_s(\xi)+\sum_{l=-\infty}^{+\infty} W_l\psi_l e^{i\lambda_l\tau}, \tag{8}$$

where $\psi_l$ is the $l$'th supermode with a complex gain coefficient $\lambda_l$ [5].

Typically, the linearly unstable supermodes $\psi_l$, which saturate nonlinearly by depleting the electron energy and inducing electron energy spread, dominate any steady state solution. This is not necessarily the case. We demonstrate below how the IRKA concept is realized by a high-$Q$ cavity operating near synchronism with *no linearly unstable* supermodes, provided that the entering electron bunches are slightly prebunched.

## III. ANALYSIS OF REGENERATIVE AMPLIFICATION

The method used to prebunch the electron beam is unimportant to the operation of the IRKA. Here we assume prebunching at the desired radiation wavelength and concentrate on the physics of pulse generation in the IRKA configuration. If a low-power long-pulse radiation source at the desired wavelength is available (as may be the case for microwave frequencies), prebunching can be achieved in a separate single-pass bunching structure which is powered by an external source. Alternatively, microbunching may be generated by sending an electron beam through a masked magnetic chicane [9]. Simple estimates, confirming the plausibility of masked chicane microbunching for a typical infrared FEL are given in Sec. IV. If the external source is not available (as may be the case for infrared frequencies), and masked chicane microbunching is not practical (e.g., because the energy slew required for microbunching is excessive for lasing), a separate oscillator may be needed to prebunch the beam, as shown in Fig. 1. The prebuncher must be *linearly unstable* and operate as a low-power FEL oscillator. The initial radiation in the prebuncher is seeded by an infinitesimally small shot noise of beam electrons. However, radiation grows over time and saturates nonlinearly, reaching the steady state. In steady state, electrons interact with radiation and develop a finite bunching level toward the end of the wiggler (as shown schematically in Fig. 1).

The IRKA configuration is governed by experimentally controlled FEL oscillator parameters, the cavity desynchronism and cavity losses, such that (i) the oscillator is operated very close to synchronism, $\nu\ll 0.13$, and (ii) there are no linearly unstable supermodes. $\nu=0.13$ corresponds to the lasing threshold for a lossless oscillator. Condition (ii) implies

$$\alpha\gg 3\sqrt{3}(\nu/2)^{2/3}. \tag{9}$$

Hence after a time $\tau\approx 1/\alpha$, a steady state is achieved. Equations (2)–(6), with $d/d\tau=0$, can then be used to model the steady-state operation. We observe from numerical simulations that ultrashort radiation pulses are generated in the linear regime, where $Q$ and $S$ can be assumed constant. The linearized steady-state equations can then be recast as

$$-\nu\ \frac{\partial A}{\partial \xi}+\frac{\alpha}{2}A=\eta B_1 \tag{10}$$

and

$$\left(\frac{\partial^2}{\partial \xi^2}+\sigma^2\right)(B_1 e^{iy_0\xi})=iAe^{iy_0\xi}, \tag{11}$$

where we assumed $Q=y_0$, and $S=y_0^2+\sigma^2$. Equation (10) describes the steady-state radiation emitted by a prebunched electron beam and deserves a careful physical interpretation.

Even though Eq. (10) is independent of time, it contains the physics associated with laser lethargy. For example, radiation emitted by the microbunched electron beam not too close to the wiggler exit experiences a strong lethargy, much exceeding the effect of cavity desynchronism. This lethargy pushes radiation toward $\zeta=1$ as it amplifies. Thus, for $\zeta<\zeta_{\max}$, where $\zeta_{\max}$ is close to 1 for small cavity desynchronism $\nu\ll 0.13$, the first term in the left-hand side of Eq. (10) can be neglected. This amounts to assuming a perfectly synchronized cavity. The analysis of the transient behavior of radiation in a perfectly synchronized cavity was carried out by Piovella [10], and will be shown to be in qualitative agreement with our results. In contrast, in the immediate vicinity of $\xi=1$ cavity desynchronism dominates over the laser lethargy, and the radiation emitted by a highly bunched electron beam at the exit from the wiggler, is pushed toward $\xi=0$ by the finite cavity desynchronism. The buildup of $|A|$ from zero at $\xi=1$ to its peak value is mainly governed by an interplay of the strong bunching [the right-hand side of Eq. (10)] and the cavity desynchronism.

Equation (10) is a typical example of a boundary layer problem, where a term with the highest (first) derivative is multiplied by a small parameter $\nu$. Hence Eqs. (10) and (11) can be solved between $\xi=0$ and $\xi=\xi_{\max}\approx 1$ by neglecting the effect of cavity desynchronism, and then between $\xi=\xi_{\max}$ and $\xi=1$ to insure the boundary condition $A(\xi=1)=0$. This approach will be justified by the results of the calculation and the assumption of a strongly damped cavity. Note that *lengthening* the cavity ($\nu<0$) is counterproductive to generation of ultrashort pulses. This is because the peak value of $|A|$ will occur at $\xi=1$, and a significant pedestal will be formed, due to the high $Q$ of the cavity, for $\xi>1$. By operating the optical resonator with a *shortened* cavity ($\nu>0$), the peak value of $|A|$ is shown below to greatly exceed $|A|(\xi=0)$, making the pedestal formation for $\xi<0$ insignificant.

To illustrate our approach to this boundary layer problem, assume, for simplicity, $\nu y_0/\alpha \ll 1$. An approximate solution to Eqs. (10) and (11) is then

$$A(\xi)=A_0(\xi)-A_0(\xi=1)e^{\alpha(\xi-1)/2\nu}, \tag{12}$$

where $A_0(\xi)$ and $B_0(\xi)$ are solutions far away from the boundary layer at $\xi=1$, and

$$A_0(\xi)\approx \frac{2B_0(\xi)}{\alpha}\left(1-\frac{2i\nu y_0}{\alpha}\right). \tag{13}$$

$B_0(\xi)$ in Eq. (13) is obtained by solving the linear equation

$$\left(\frac{\partial^2}{\partial\xi^2}+\sigma^2\right)B_0 e^{iy_0\xi}=i\frac{2}{\alpha}\ (1-2i\nu y_0/\alpha)B_0 e^{iy_0\xi}, \tag{14}$$

with boundary conditions $B_0(\xi=0)=b_0$ and $dB_0/d\xi(\xi=0)=0$. The largest amplitudes of the radiation field are expected when the energy spread vanishes, i.e., for $\sigma=0$. More realistic cases with $\sigma\neq 0$ are studied numerically.

The solution of Eq.(14) with $\sigma=0$ is

$$A_0(\xi)=\frac{b_0}{\alpha+2i\nu y_0}[(1+iy_0/k_1)e^{k_1\xi}+(1-iy_0/k_1)e^{-k_1\xi}], \tag{15}$$

where

$$k_1\approx\frac{1}{\sqrt{\alpha}}[(1+\nu y_0/\alpha)+i(1-\nu y_0/\alpha)]. \tag{16}$$

Equations (15) and (12) indicate that the half-width of the laser pulse is of order $\sigma_-=\sqrt{\alpha}$ for $\xi<\xi_{\max}$, and $\sigma_+=2\nu/\alpha$ for $\xi<\xi_{\max}$. Because of condition (9) for operating the IRKA under lasing threshold, $\sigma_-\gg\sigma_+$.

Analysis of the transient evolution of a perfectly synchronized FEL oscillator [10] facilitates a qualitative understanding of the amplitude and duration of the radiation pulses in the regenerative amplifier described here. Consider a single radiation pulse injected into the cavity at $\tau=0$ (which can also be accomplished by microbunching one of the entering electron bunches). In a perfectly synchronized cavity the radiation pulse evolves nonexponentially according to $|A|^2(\zeta,\tau)\propto\exp(3\sqrt{3}/2^{2/3}\zeta^{2/3}\tau^{1/3}-\alpha\tau)$. Initially, for $\tau<\tau_{\max}=1/2(\sqrt{3}/\alpha)^{3/2}$, the radiation pulse builds up, narrows, and then decays exponentially [5]. The peak amplitude of the radiation pulse $|A|^2(\zeta,\tau_{\max})\propto\exp[\zeta(3\sqrt{3}/\alpha)^{1/2}]$, and its width is $\sigma\propto\sqrt{\alpha}$. Radiation in the cavity does not decay if a train of radiation pulses (or microbunched beams) is injected. Instead, a steady state is reached, in which a radiation pulse is a sum of the transient solutions corresponding to different injection times $\tau$. At any time $\tau$ the most significant contribution comes from the pulse injected at $\tau'=\tau-\tau_{\max}$. Hence the steady-state solution has the same qualitative properties as $|A|^2(\zeta,\tau_{\max})$, i. e. its amplitude scales as $\exp(\alpha^{-1/2})$ and its width scales as $\alpha^{1/2}$ [in agreement with Eq. (15)].

From Eqs. (15) and (16) it follows that the efficiency of the IRKA increases with $y_0$, for $y_0\ll\alpha/\nu$. In other words, it is advantageous to inject pre-bunched electrons at an energy exceeding the resonant energy of the IRKA. In case of a FEL prebuncher this can be achieved by changing the wavelength and/or the number of periods of IRKA wiggler with respect to those of the prebuncher wiggler. The change in the dimensionless detuning $\Delta y_0=y_{\mathrm{IRKA}}-y_{\mathrm{bunch}}$ is given by

$$\Delta y_0=\frac{\Delta(N_w\lambda_w)}{N_w\lambda_w}y_{\mathrm{bunch}}-2\pi N_w\frac{\Delta\lambda_w}{\lambda_w}, \tag{17}$$

where $y_{\mathrm{bunch}}$, $\lambda_w$, and $N_w$ are the normalized energy detuning, wiggler wavelength, and number of wiggler periods in the prebuncher, respectively, and $y_{\mathrm{IRKA}}$, $\lambda_w+\Delta\lambda_w$, and $N_w+\Delta N_w$ are the corresponding quantities in the IRKA. Assuming that the number of wiggler periods is the same for both wigglers, Eq. (17) indicates that for $\Delta y_0=10.0$ and $N_w=50$ the fractional decrease in wiggler wavelength is about 3.2%.

## IV. EXAMPLE

Below we demonstrate how radiation pulses much shorter than a slippage length can be generated by a typical infrared FEL, such as compact infrared FEL (CIRFEL) [4]. For typical parameters of $Q=1.5$ nC, $a_w=0.25$, $N_w=50$, $\lambda_w=1.0$ cm, $r_b=1$ mm, $E=8$ MeV, and $L_b\approx 1$ mm, the resonant wavelength is $\lambda_r=20$ $\mu$m, the slippage length is approximately equal to bunch length, and the gain parameter $\Gamma\approx 4$. By choosing cavity loss $\alpha_0=20\%$, cavity desynchronism $\delta\mathcal{L}=2\mu$, energy spread $\delta\gamma/\gamma=0.5\%$, and assuming that electrons are microbunched with periodicity $\lambda=20.6\mu$, the normalized parameters are calculated to be $y_0=10.0$, $\sigma=1.2$, $\alpha=0.05$, and $\nu=0.001$. Results of a time-dependent numerical simulation of the nonlinear Eqs. (2)–(6), with initial prebunching $b_0=0.0025$, are presented in Fig. 2. The full width at half maximum of the optical pulse in Fig. 2 is about one-sixth of the slippage length.

Micro-bunching on a $\lambda=20.6\mu$ scale using a masked chicane appears to be technologically feasible on CIRFEL. Following the calculation in Ref. [9], consider a four-magnet chicane with a thin-wire mask in the middle. If an energy slew is imposed on a beam by accelerating electrons on the sloped part of the rf bucket, a correlation exists between

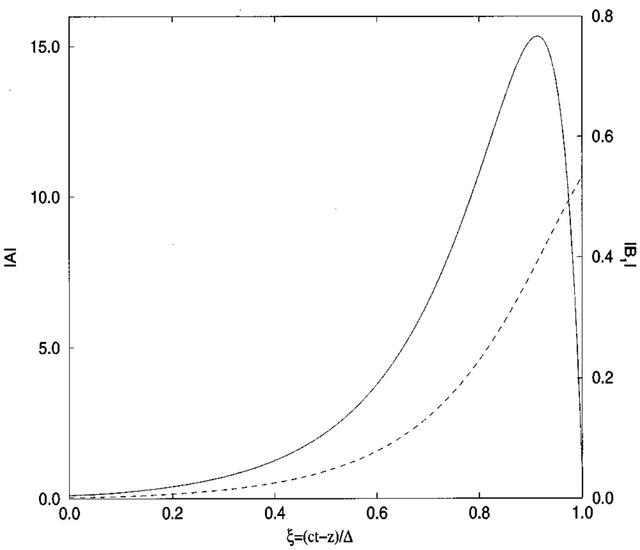


FIG. 2. Normalized field amplitude and bunching in IRKA (one slippage window) for typical CIRFEL parameters: $Q=1.5$ nC, $a_w=0.25$, $N_w=50$, $\lambda_w=1.0$ cm, $r_b=1$mm, $E=8$ MeV, $L_b\approx 1$ mm, $\lambda_r=20$ $\mu$m, $\lambda=20.6\mu$, $\alpha_0=20\%$, $\delta\mathcal{L}=2\mu$, and energy spread $\delta\gamma/\gamma=0.5\%$. Solid line, field amplitude; dashed line, bunching. The results are a numerical solution to Eqs. (2)–(6).

energy and longitudinal position. The chicane can be designed to be longitudinally achromatic. However, in the middle of the chicane electrons with different energies become dispersed transversely. A series of thin wires, spaced by distance $d$ in the transverse direction, periodically perturbs the beam by knocking out some of the electrons. This perturbation translates into a longitudinal density perturbation behind the chicane. For a chicane consisting of four bending magnets of length $L=30$ cm with magnetic field $B=5$ kG, separated by drifts $D=10$ cm, and assuming the beam energy slew of 0.5%, we find that the longitudinal perturbation is produced with periodicity $\lambda\approx 0.14d$. Choosing the mask periodicity $d=150$ $\mu$m produces the required bunching with $\lambda=20.6$ $\mu$m. Bunching periodicity can be tuned by adjusting the magnetic field of the chicane.

An important point to have in mind is that narrow single-spike solutions, as seen in Fig. 2, can only be obtained in the linear regime, where overbunching of the electrons does not lead to nonlinear saturation. To insure that the IRKA operates in a linear regime, initial prebunching must be small (0.25% in the numerical example shown in Fig. 2). The magnitude of microbunching may be adjusted by propagating the beam through a dispersive section which debunches the beam. Energy slew, introduced in order for the masked chicane prebunching to work, will not severely affect the FEL performance in the IRKA section as long as $2\pi N_w\delta\gamma/\gamma\ll\sqrt{2/\alpha}$. This is because radiation pulses of width $\sigma_-=\sqrt{\alpha}\Delta$ are very narrow, making them insensitive to energy spread.

## V. CONCLUSIONS

In conclusion, we have suggested a FEL configuration: the injection-locked regenerative klystron amplifier (IRKA). The IRKA configuration offers a way of producing a steady-state train of ultrashort radiation pulses, while avoiding the chaotic dynamics characteristic of operating a FEL oscillator near cavity synchronism. Our analysis has been restricted to bunch lengths much less than a slippage length. We conjecture that a similar multicavity scheme would work for electron bunches longer than the slippage length. As the cavity desynchronism decreases, the radiation pulse is pushed off the electron beam by lethargy, and is narrowed to a spike, which trails the electron bunch. Future work will remove this assumption and examine the regenerative amplification regime for a wide variety of FEL systems. Extending the one-dimensional analysis to include laser diffraction and beam emittance requires further investigations.

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[1] G. M. H. Knippels, R. F. X. A. M. Mols, A. F. G. van der Meer, D. Oepts, and P. W. van Amersfoort, Phys. Rev. Lett. **75**, 1755 (1995).

[2] H. A. Schwettman, Nucl. Instrum. Methods Phys. Res. A **375**, 156 (1996).

[3] D. A. Jaroszynski *et al.* Phys. Rev. Lett. **71**, 3798 (1993).

[4] I. S. Lehrman, J. Krishnaswamy, R. A. Hartley, R. H. Austin, and D. W. Feldman, Nucl. Instrum. Methods Phys. Res. Sect. A (to be published).

[5] N. Piovella, P. Chaix, G. Shvets, and D. A. Jaroszynski, Phys. Rev. E **52**, 5470 (1995); N. Piovella, P. Chaix, D. A. Jaroszynski, and G. Shvets, Nucl. Instrum. Methods Phys. Res. Sect. A **375**, 156 (1996).

[6] G. Dattoli, A. Marino, F. Romanelli, Opt. Commun. **35**, 407 (1980); G. Dattoli, A. Marino, A. Renieri and F. Romanelli, IEEE J. Quantum Electron. **17**, 1371 (1981).

[7] N. S. Ginzburg, S. P. Kuznetsov, and T. N. Fedoseeva, Radiophys. Quantum Electron. **21**, 728 (1971).

[8] R. Bonifacio, F. Casagrande, and L. De Salvo Souza, Nucl. Instrum. Methods Phys. Res. Sect. A **33**, 2836 (1986); R. Bonifacio, L. De Salvo Souza, P. Pierini, and N. Piovella, *ibid.* **296**, 358 (1990).

[9] D. C. Nguyen and B. E. Carlsten, Nucl. Instrum. Methods Phys. Res. Sect. A **375**, 597 (1996).

[10] N. Piovella, Phys. Rev. E **51**, 5147 (1995).