PHASE-STABLE, MICROWAVE FEL AMPLIFIER*

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Free-electron laser (FEL) amplifiers have demonstrated high efficiency and high output power for microwave wavelengths. However, using present technology, microwave FEL amplifiers are not phase stable enough to be suitable for driving linear accelerators, where several such amplifiers need to be phase locked. The growing wave's phase sensitivity to the beam voltage in the small-signal regime is responsible for the largest contribution to this phase instability. We discuss a scheme that reduces the phase sensitivity to the beam voltage by operating off synchronism and matching the phase variation resulting from the desynchronism to the phase variation from the reduced plasma wavenumber as the beam voltage changes.

I. INTRODUCTION

Free-electron lasers (FELs) have demonstrated both high beam-to-rf power extraction efficiencies (~30%) [1] and high output power (on the order of gigawatts) [2], and have been considered as candidates to drive high-frequency advanced accelerators like those proposed for linear colliders [3]. However poor phase stability has been predicted and measured for FELs [4,5]. Typical accelerator applications require rf phase stability on the order of 5^{o} of phase, and advanced accelerator applications such as bunch compression [6] and short-wavelength FELs require stability to 1^{o} or less [7]. Phase noise in microwave FELs arises from fluctuations in tube voltage, current, confining magnetic field strength, and other tube parameters. Typically, the largest effect is from voltage fluctuations. Electron beams for practical FELs used as rf sources will have diode voltages of 1/2 to 1 MV with voltage stabilities on the order of 1/4%. Measured and simulated FEL phase stability to date has been on the order of a 20^{o} to 40^{o} shift per percent voltage fluctuation [5,8,9]. This level of phase stability does not satisfy advanced accelerator requirements.

In a klystron, the phase of a cavity is completely determined by the absolute phase of the harmonic current at that location. If the beam energy is shifted slightly by $\delta\gamma$, we can expect that the output phase will shift by $\delta\Phi = -(\beta_e L / \gamma(\gamma+1))(\delta\gamma / (\gamma-1))$ where L is the total device length and the electron propagation number is $\beta_e = \omega / \beta c$, and where ω is the frequency of operation and β is the beam axial velocity normalized to the speed of light. For a 500 keV beam in a half-meter-long 11.4 GHz klystron, a 1/4% shift in the voltage will lead to about a 3^{o} phase shift. Note that if the product ωL is kept constant, tubes at other frequencies will have the same phase shift for the same voltage shift. However, since the space-charge bunching length is independent of operating frequency, the product ωL will in general increase as the frequency is increased and we can expect that a 20 GHz tube will have phase variations on the order of 5^{o}, and higher frequency tubes will have even larger phase variations. Lower frequency klystrons typically have phase stability better than 1^{o} [7].

The phase shift in an FEL due to the transit time effect is the same as for a klystron. However, the effect of the space-charge wave and the transit time of the electron beam are not separable in a FEL as they are in a klystron. This will introduce new physical effects, one of which is the possibility of using fluctuations in the space-charge wave to counter fluctuations in the beam's transit time through the device.

In this paper we will explicitly demonstrate that the phase dependency on the space-charge wave can effectively cancel the phase dependency on the beam's transit time factor for the FEL interaction. We will do this by analyzing the dispersion relation for an axial FEL, which is used instead of the conventional transverse FEL for simplicity. We will also present numerical solutions of the dispersion relation exhibiting the phase-stable condition.

II. DISPERSION RELATION FOR AN AXIAL FEL

Recently, an axial FEL interaction was proposed for the generation of gigawatt microwave radiation [10]. In this device, an annular electron beam interacts with the axial electric field of a TM_{0m} mode in a circular waveguide. The radius of this waveguide is periodically rippled which causes the mode to radially expand and contract. The ripple amplitude is only a few percent of the average radius, and the mode is able to adiabatically conform to the gradual change in the waveguide radius. The axial FEL interaction for a synchronous particle is shown in Figure 1. The annulus is located at a radius corresponding to a zero of the axial electric field of that mode in a waveguide with a radius equaling the mean radius of the rippled waveguide. When an electron is at the axial position of the smallest waveguide radius the axial electric field at the location of the electron opposes the electron's motion. As the electron travels to the region of larger radius the rf slips by the electron. When the electron is at the location of the maximum waveguide radius one half of a rf wavelength has slipped by, resulting in a sign change in the mode's fields. Additionally, the electron is experiencing the electric field at a radius larger than the axial field zero instead of a radius smaller. This switch from one side of the null of the axial electric field to the other provides another sign change in the axial field at the location of the electron, and the electric field is again opposing the electron's motion. This interaction is equivalent to the interaction of a

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transverse-coupling FEL except the rf field is wiggled instead of the electrons to provide synchronism.

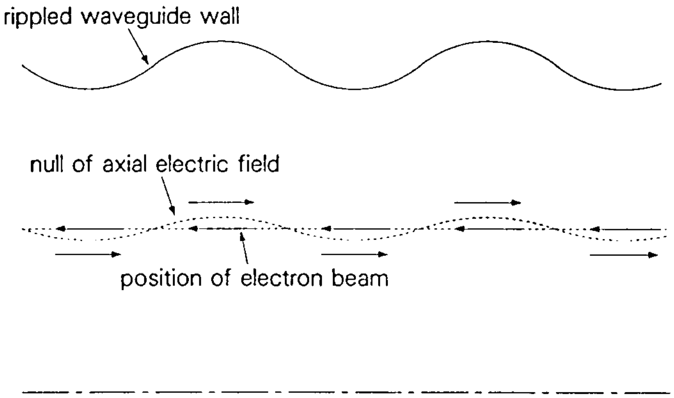

Figure 1: Axial electric field orientations for a synchronous particle when the particle reaches the centers of the ripples in an axial FEL, in *r-z* geometry

If we assume that the growing rf mode components have an $e^{j\omega t-\Gamma z}$ exponential behavior, this dispersion relation can be derived [11] for the axial FEL:

$$\left((\beta_e + j\Gamma)^2 - \left((j\Gamma)^2 - k^2\right)\hat{\beta}_q^2\right)\left((\Gamma - jk_w)^2 + \beta_1^2\right) = 2C^3\beta_1^4 \tag{1}$$

where the cold rf mode propagation constant is $\beta_1 = \omega / v_{phase}$, the gain parameter C is proportional to the beam current divided by $\gamma\beta^3$, all to the 1/3 power, and the normalized space-charge wavenumber is $\hat{\beta}_q^2 = 4\chi_o \frac{I}{I_A} \ln \frac{r_w}{r_b}$, where now r_w is the wall radius, r_b is the beam annulus radius, I is the beam current, I_A is about 17 kA, and χ_o is a slowly varying term close to unity that depends on r_w and r_b. We can define the normalized gain δ and the detuning Δ by

$$\begin{aligned} \Gamma &= j\beta_1 + jk_w + \delta C\beta_1 \\ \beta_e &= \beta_1 + k_w + \Delta \end{aligned} \tag{2}$$

where k_w is the wiggler wavenumber, given by 2π divided by the wiggler period. Eq. (1) now becomes

$$\begin{aligned} &\Big(\Delta^2 + 2j\Delta\delta C\beta_1 - (\delta C\beta_1)^2 \\ &-\left((\beta_1 + k_w)^2 - k^2 - 2j\delta C\beta_1(\beta_1 + k_w)\right)\hat{\beta}_q^2\Big)(j\delta) = C^2\beta_1^2 \quad . \end{aligned} \tag{3}$$

Now δ and thus the output phase can be make stable if

$$\begin{aligned} \frac{d}{d\beta_e}\Big(\Delta^2 + 2j\Delta C\beta_1 - \big((\beta_1 + k_w)^2 - k^2 \\ - 2j\delta C\beta_1(\beta_1 + k_w)\big)\hat{\beta}_q^2\Big) = 0 \quad , \end{aligned} \tag{4}$$

for the case the derivative of the interaction strength with respect to β_e vanishes. Eq. (4) is satisfied by the conditions (for a constant perveance gun):

$$\begin{aligned} \Delta &= -\beta_e / \gamma \\ \hat{\beta}_q^2 &= \frac{2}{3\gamma(\gamma+1)(1-\Delta/\beta_e)} \end{aligned} \quad . \tag{5}$$

This solution only makes sense for γ on the order of 10 or greater because of the typically narrow window of detunings that lead to growing mode solutions. However, we will next show that the first equation in Eq. (5) is required for gain stability, and a minor modification to the second equation will lead to phase stability for small detunings.

Let us assume that $\Gamma = j\beta_1 + jk_w + \delta_o C\beta_1$ is a solution of the dispersion relation, Eq. (1). Now let us consider the solution of the dispersion relation where Δ is slightly shifted (by δ_Δ), $\hat{\beta}_q^2$ is slightly shifted (by $\delta_{\hat{\beta}_q^2}$), and C is slightly shifted (by δ_C), and where we denote the new solution by $\Gamma = j\beta_1 + jk_w + \delta_o C\beta_1 + \delta_1 C\beta_1$. Since the solution to the growing mode has a negative real component and lags behind the electron's phase velocity, we can write $\delta_o = -a + jb$, where both a and b are positive and typically on the order of unity. After solving for δ_1 by performing a first-order expansion we find

$$\hat{\beta}_q^2 = \frac{2}{3\gamma(\gamma+1)} \times \left(\frac{1 - \frac{b\Delta}{3(a^2+b^2)C\beta_1} - \frac{dC/d\gamma}{C}\frac{3(a^2-b^2)C\beta_1\beta^2\gamma^3}{2(a^2+b^2)^2\beta_e}}{(1-\Delta/\beta_e) + \frac{b\left((\beta_1+k_w)^2 - k^2\right)}{2(a^2+b^2)C\beta_1\beta_e}} \right) \tag{6}$$

as the condition for phase-stable operation.

III. NUMERICAL SOLUTIONS OF THE DISPERSION RELATION SHOWING PHASE STABILITY

We can numerically find the growing root of the dispersion relation, Eq. (1). In this section we will do this for a variety of beam energies, demonstrating the phase-stable conditions found in the last section over a broad operating range.

A. High energy, low gain

For the case of high energy and low gain, the solution specified in Eq. (5) is valid if the interaction strength is independent of beam energy. For the case $\gamma = 100$, $\beta_e = 300$ m^{-1}, and $C = 0.03$, this solution is given by $\Delta = -3$ m^{-1} and $\beta_q^2 = 6.7(10^{-5})$. In Figure 2 we plot the derivatives of the phase change per unit length and the amplitude growth with respect to beam energy, respectively, as calculated numerically from Eq. (1) for $\Delta = -3\ \mathrm{m}^{-1}$ while varying β_q^2, and while assuming the interaction strength is independent of beam energy and the beam has constant perveance. As predicted, both derivatives vanish at $\beta_q^2 = 6.7(10^{-5})$, which is an autostable operating point.

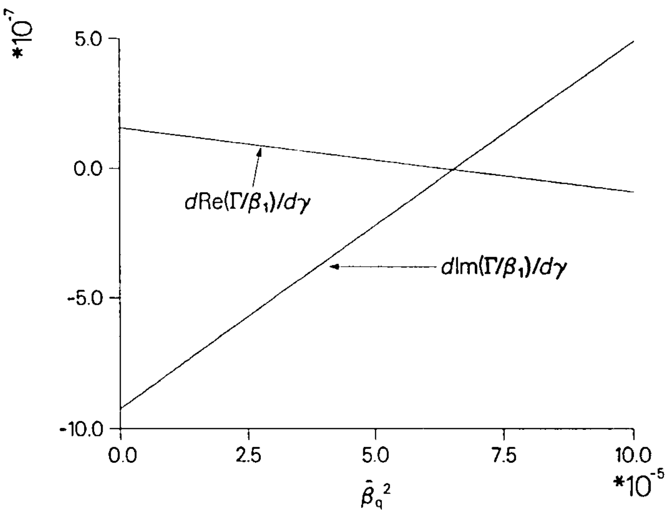

Figure 2: Sensitivity of phase and gain to beam energy for low gain, high energy case satisfying Eq. (5) as a function of space-charge wavenumber.

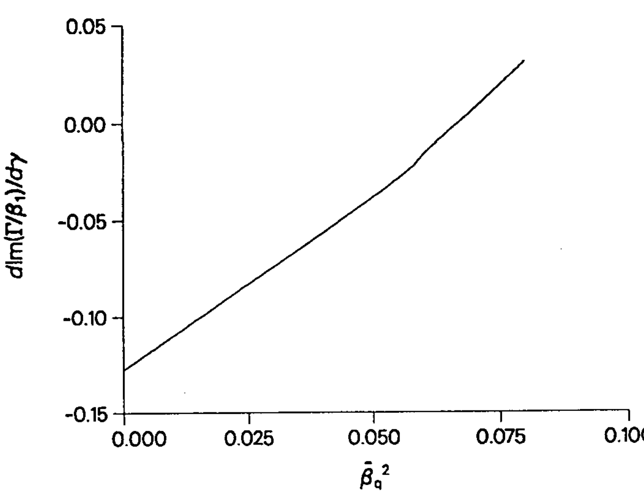

Figure 3: Sensitivity of phase to beam energy for medium gain ($C = 0.1$), low energy ($\gamma = 2$) case as a function of the space-charge wavenumber, for a detuning $\Delta = -50\ \mathrm{m}^{-1}$.

B. Low energy, moderate gain

Now let us consider another constant perveance case with $\gamma = 2$ at 13 GHz (so the beam propagation constant is about 300 m^{-1}), an output power of about 1 GW, and with a device length of about 1 m. For these parameters the gain constant C is on the order of 0.1. For a detuning of $\Delta = -50\ \mathrm{m}^{-1}$, Eq. (6) predicts phase-stable operation at a space-charge wave number of about 0.08. In Figure 3 we have plotted the derivative of the phase change per unit length with respect to beam energy as a function of the space-charge wave number numerically calculated for this detuning. The calculated growth rate is about 11 m^{-1}, or about a factor of 2 per wiggler period if the wiggle wave number $k_w = 100\ \mathrm{m}^{-1}$(a 6 cm wiggler period), and phase-stable operation is achieved with a beam current of about 5 kA at nearly the predicted space-charge wave number.

IV. CONCLUSION

We have examined the FEL dispersion relation in the Raman regimes and shown that if the interaction strength is independent of beam energy, there is a stable phase and gain operating point. We have additionally shown that even if the interaction strength depends on the beam energy there is a phase-stable operating point, which we then demonstrated numerically. It should be noted that this technique of reducing the phase stability of an FEL is not possible with a klystron, and that by proper design an FEL can have a phase stability an order of magnitude greater (or more) than a klystron.

V. REFERENCES

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