

Analytical analysis of longitudinal space charge effects for a bunched beam with radial dependence *

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The longitudinal space-charge (LSC) force can be a major cause of the microbunching instability in the linac for an x-ray free-electron laser. In this paper, the LSC-induced beam modulation is studied using an integral equation approach that takes into account the transverse (radial) variation of the LSC field for both the coasting-beam limit and a bunched beam. Variation of the beam energy and the transverse beam size is also incorporated. We discuss the validity of this approach and compare it with other analytical analyses as well as numerical simulations.

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I. INTRODUCTION

To ensure the successful commissioning and operation of an x-ray free-electron laser (FEL), the highest quality electron beam is prepared [1]. However, such a high-quality electron beam is subject to various instabilities along the accelerator system. Indeed, the FEL is one of the collective instabilities happening in the undulator. Because of the fact that the electron beam born in the radio frequency (rf) gun has very small energy spread, Landau damping is typically ineffective in suppressing all of the unwanted instabilities [2–4]. Because of inevitable density granularity of the electron beam from the rf cathode, the longitudinal space-charge (LSC) effect can induce a large energy modulation on the beam, that leads to instability downstream when the beam is accelerated and compressed [3, 5]. In this paper, we study the LSC effect, taking into account acceleration and the variation of transverse beam size during the acceleration. We study this analytically via an integral equation approach, which is then compared to direct numerical simulations, and also to other analytical approaches [6].

The paper is organized as follows. In Sec. II, we introduce a coasting-beam theory. We first work out details for a one-dimensional (1D) theory, and then extend the theory to the case where the impedance has a radial dependence. In Sec. III, we advance the theory to study a bunched-beam case, where a multifrequency theory is developed. Again, both a 1D theory and a theory for the impedance with a radial dependence are developed. In Sec. IV, we apply this general framework to the case where the impedance comes from the LSC effect. The LSC impedance of a δ -ring with a radius r is presented.

With this r -dependent impedance, we provide some examples to compare our results with other analytical theories and numerical simulation in Sec. V. We conclude with a discussion in Sec. VI.

II. COASTING-BEAM THEORY

We first study the case where the electron beam is approximated as a coasting beam. This approximation has been adopted in previous work [2–4, 6, 7]. For a coasting-beam case, the study gives a single-frequency theory. We derive integral representations for the density modulation and the energy modulation for the case where the electron beam is being accelerated, and its transverse size varies along the beam line.

A. One-dimensional formulae

We describe the electron beam by a distribution function $f(x, x', y, y', z, \gamma; s)$. Here, x , y , and z are the internal coordinates in the electron beam; x' and y' are the slopes of the trajectories; γ is the electron total energy in units of the electron rest energy mc^2 with m the electron mass and c the speed of light in vacuum; and $s = c \int \sqrt{1 - \gamma^{-2}} dt$ is the position along the beam line. The distribution function immediately after the energy kick due to the wakefield (at $\tau + 0$) is related to that immediately before (at $\tau - 0$) by

$$\begin{aligned} f(\mathbf{X}_\tau; \tau + 0) &= f(\mathbf{X}_\tau - \Delta\mathbf{X}; \tau - 0) \\ &\approx f(\mathbf{X}_\tau; \tau - 0) - \Delta\gamma \frac{\partial f(\mathbf{X}_\tau; \tau - 0)}{\partial \gamma}; \end{aligned} \quad (1)$$

where $\mathbf{X} \equiv (x, x', y, y', z, \gamma)$ for the six-dimensional phase space coordinates, and $\Delta\mathbf{X} = (0, 0, 0, 0, 0, \Delta\gamma)$. Here, we focus on the longitudinal phase space only. Summing up the wakefield contribution over the entire trajectory, (*i.e.*, $\tau \in [0, s]$), and using the boundary notations $f[\mathbf{X}_{\tau \rightarrow s}; (\tau \rightarrow s) + 0] = f[\mathbf{X}(s); s]$ and $f[\mathbf{X}_{\tau \rightarrow 0}; (\tau \rightarrow$

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0) - 0] = $f_0(\mathbf{X}_0)$, the evolution of the distribution function under the influence of the wakefield is

$$f[\mathbf{X}(s); s] = f_0(\mathbf{X}_0) - \int_0^s d\tau \frac{\partial f(\mathbf{X}_\tau; \tau - 0)}{\partial \gamma_\tau} \frac{d\gamma}{d\tau}. \quad (2)$$

The rate of energy change due to the wakefield is

$$\frac{d\gamma}{d\tau} = -r_e \int \frac{dk_1}{2\pi} Z(k_1; \tau) N b(k_1; \tau) e^{ik_1 z_\tau}. \quad (3)$$

In Eq. (3), r_e is the electron classical radius, N is the total number of electrons, $Z(k_1; s)$ is the longitudinal impedance, and we have introduced the density bunching factor $b(k; s)$ as

$$b(k; s) = \frac{1}{N} \int d\mathbf{X} e^{-ikz} f(\mathbf{X}; s), \quad (4)$$

where k is the wavenumber characterizing the frequency dependence of the density bunching factor. Using Eq. (2) and the density bunching factor defined in Eq. (4), we have

$$b[k(s); s] = b_0[k(s); s] - \frac{ik}{N} \int d\tau \mathcal{R}_{56}(\tau \rightarrow s) \times \int d\mathbf{X}_\tau e^{-ikz(\mathbf{X}_\tau)} f(\mathbf{X}_\tau; \tau - 0) \frac{d\gamma}{d\tau}, \quad (5)$$

where

$$b_0(k; s) = \frac{1}{N} \int d\mathbf{X}_0 e^{-ikz} f_0(\mathbf{X}_0) \quad (6)$$

is the bunching factor without wakefield. We have introduced the transfer matrix as $\mathbf{X}(s) \equiv \mathcal{R}(\tau \rightarrow s)\mathbf{X}(\tau)$. Plugging Eq. (3) into Eq. (5), we have

$$b[k(s); s] = b_0[k(s); s] + ik r_e \int d\tau \mathcal{R}_{56}(\tau \rightarrow s) \times \int \frac{dk_1}{2\pi} Z(k_1; \tau) b(k_1; \tau) \int d\mathbf{X}_0 e^{-ikz + ik_1 z_\tau} f_0(\mathbf{X}_0), \quad (7)$$

where $z_\tau = z_0 + \mathcal{R}_{56}(0 \rightarrow \tau)\Delta\gamma_0$, and $z = z_0 + \mathcal{R}_{56}(0 \rightarrow s)\Delta\gamma_0$.

We assume that the initial distribution function can be decomposed into two parts

$$f_0(\mathbf{X}_0) = \bar{f}_0(\mathbf{X}_0) + \hat{f}_0(\mathbf{X}_0), \quad (8)$$

where $\bar{f}_0(\mathbf{X}_0)$ is the average distribution function, and $\hat{f}_0(\mathbf{X}_0)$ is the initial microbunching. For microbunching wavelengths much smaller than the electron beam length, we can assume a uniform longitudinal distribution, hence a coasting beam in z , and Gaussian in $\Delta\gamma$ for the average distribution function, in other words, we assume

$$\bar{f}_0(\mathbf{X}_0) = \frac{n_0}{\sqrt{2\pi}\sigma_{\Delta\gamma}} \exp\left\{-\frac{\Delta\gamma^2}{2\sigma_{\Delta\gamma}^2}\right\}, \quad (9)$$

where $n_0 = N/L$ is the average line density with L being the electron beam length. Within the linear theory, we

can neglect $\hat{f}_0(\mathbf{X}_0)$ in completing the integral in Eq. (7). In doing so, we get

$$b[k(s); s] = b_0[k(s); s] + \int_0^s d\tau K(\tau, s) b[k(\tau); \tau], \quad (10)$$

with the kernel of the integral equation as

$$K(\tau, s) = ik(s)\mathcal{R}_{56}(\tau \rightarrow s) \frac{I(\tau)Z[k(\tau); \tau]}{I_A} \times \exp\{-k_0^2 \mathcal{R}_{56}^2(\tau \rightarrow s)\sigma_{\Delta\gamma}^2/2\}, \quad (11)$$

where

$$\mathcal{R}_{56}(\tau \rightarrow s) \equiv \int_\tau^s \frac{dx}{\gamma(x)^3 [1 - \gamma(x)^{-2}]}. \quad (12)$$

It is worth emphasizing that, for a uniform average distribution in z and single-frequency initial microbunching, there is only a single frequency selected in the k -integral in Eq. (7). When we work on a Gaussian distribution in z , we deal with a multifrequency theory, which will be explored later.

According to Eq. (3), the resulting accumulated energy modulation spectrum is

$$\Delta\gamma[k(s); s] = - \int_0^s d\tau \frac{I_0 Z[k(\tau); \tau] b[k(\tau); \tau]}{I_A} \times \exp\{-k^2 \mathcal{R}_{56}^2(\tau \rightarrow s)\sigma_{\Delta\gamma}^2/2\}, \quad (13)$$

where $I_0 = ecn_0$ is the peak current with e the electron charge, and $I_A \approx 17045$ Amp the Alfvén current.

B. Radial Dependence

If the transverse dynamics and the longitudinal dynamics are separable, we can assume that the distribution function is factorable

$$f(\mathbf{X}; s) = f_r(r; s) f_z(z; s), \quad (14)$$

with the normalization of

$$\int d\mathbf{r} f_r(r; s) = 1. \quad (15)$$

Then the three-dimensional (3D) problem can be simplified to a one-dimensional problem.

1. Transverse averaging approach

One approach is to average out the transverse variables. The energy change rate is then

$$\begin{aligned} \frac{d\gamma(r, z; s)}{ds} &= -r_e \int dz' dr' w(z - z', r, r') f(z', r'; s) \\ &= -r_e \int dr' \int \frac{dk}{2\pi} Z(k; r, r') e^{ikz} \\ &\times \int dz' e^{-ikz'} f_r(r'; s) f_z(z'; s) \\ &= -r_e \int \frac{dk}{2\pi} \bar{Z}[k(s); r, s] N b[k(s); s] e^{ikz}, \end{aligned} \quad (16)$$

where the wakefield is introduced as

$$w(z, r, r', s) = \int \frac{dk}{2\pi} Z(k; r, r', s) e^{ikz} \quad (17)$$

and the averaged impedance is defined as

$$\bar{Z}(k; r, s) = \int d\mathbf{r}' Z(k; r, r', s) f_r(r'; s). \quad (18)$$

The bunching factor is simplified as

$$b(k; s) = \frac{1}{N} \int d\mathbf{X} e^{-ikzf}(\mathbf{X}; s) = \frac{1}{N} \int dz e^{-ikzf_z(z; s)}, \quad (19)$$

according to Eqs. (14) and (15).

Plugging Eq. (16) into Eq. (5), we have

$$b[k(s); s] = b_0[k(s); s] + ik r_e \int d\tau \mathcal{R}_{56}(\tau \rightarrow s) \int d\mathbf{X}_0 \quad (20) \\ \times \int \frac{dk_1}{2\pi} \bar{Z}(k_1; r, \tau) b(k_1; \tau) e^{-ikz + ik_1 z \tau} f_0(\mathbf{X}_0).$$

Linearizing the system and completing the integrals as in the 1D case, we formally get the same equation for the evolution of the bunching factor as in Eq. (10). However, the kernel of the integral equation is different from that given in Eq. (11). Here, the kernel is

$$K(\tau, s) = ik(s) \mathcal{R}_{56}(\tau \rightarrow s) \frac{I(\tau) \bar{\bar{Z}}[k(\tau); \tau]}{I_A} \\ \times \exp\{-k_0^2 \mathcal{R}_{56}^2(\tau \rightarrow s) \sigma_{\Delta\gamma}^2 / 2\}, \quad (21)$$

where the twice-averaged impedance is defined as

$$\bar{\bar{Z}}(k; s) = \int d\mathbf{r} \bar{Z}(k; r, s) f_r(r; s). \quad (22)$$

According to Eq. (16), the resulting accumulated energy modulation spectrum is

$$\Delta\gamma[k(s); s] = - \int_0^s d\tau \frac{I_0 \bar{\bar{Z}}[k(\tau); \tau] b[k(\tau); \tau]}{I_A} \\ \times \exp\{-k^2 \mathcal{R}_{56}^2(\tau \rightarrow s) \sigma_{\Delta\gamma}^2 / 2\}, \quad (23)$$

which should be compared to Eq. (13) for the 1D case.

2. Radial variable as parameter

Should we not perform the average in Eq. (18), we can keep the r -dependence. A similar treatment was recently adopted in Refs. [6, 7]. In our approach, we introduce a radially dependent bunching factor

$$b(k; s, r) = \frac{1}{N} \int dz e^{-ikz} f(\mathbf{r}, z; s) \\ = f_r(r; s) \frac{1}{N} \int dz e^{-ikz} f_z(z; s), \quad (24)$$

so that

$$b(k; s) = \frac{\iint_{\Sigma_{\perp}} d\mathbf{r} b(k; s, r)}{\Sigma_{\perp}}, \quad (25)$$

where $\Sigma_{\perp} \equiv \iint d\mathbf{r}$ is the transverse area. Assuming that the system has cylindrical symmetry, the 3D problem is reduced to a two-dimensional (2D) problem.

It is worthwhile to point out that in Refs. [6, 7], the r -dependence comes in as a parameter, *i.e.*, $f(\mathbf{X}; s) = f_r(r) f_z(z; s)$, but not $f_r(r; s)$ as in Eq. (14) of our paper. Without the s -dependence, $f_r(r)$ describes a constant transverse beam size.

The energy changes at the rate

$$\frac{d\gamma(r, z; s)}{ds} = -r_e \int dz' d\mathbf{r}' w(z - z', r, r') f(z', r'; s) \\ = -r_e \int d\mathbf{r}' \int \frac{dk}{2\pi} Z(k; r, r') e^{ikz} \\ \times \int dz' e^{-ikz'} f_r(r'; s) f_z(z'; s) \quad (26) \\ = -r_e \int d\mathbf{r}' \int \frac{dk}{2\pi} Z[k(s); r, r'] N b[k(s); s, r'] e^{ikz}.$$

Therefore, the bunching factor evolves as

$$b[k(s); s, r] = b_0[k(s); s, r] \quad (27) \\ + \int_0^s d\tau \int d\mathbf{r}' K(\tau, s, r, r') b[k(\tau); \tau, r']$$

with

$$K(\tau, s, r, r') = ik(s) \mathcal{R}_{56}(\tau \rightarrow s) \frac{I(\tau) Z[k(\tau); \tau, r, r']}{I_A \Sigma_{\perp}}. \quad (28)$$

The corresponding evolution for the energy modulation is then

$$\Delta\gamma(s, r) = - \int_0^s d\tau \int d\mathbf{r}' \frac{I_0 Z[k(\tau); \tau, r, r'] b[k(\tau); \tau, r']}{I_A \Sigma_{\perp}}. \quad (29)$$

Hence, the average energy modulation is

$$\Delta\gamma(s) = \frac{\iint_{\Sigma_{\perp}} d\mathbf{r} \Delta\gamma(s, r)}{\Sigma_{\perp}}. \quad (30)$$

Depending on the details of the radial dependence of $f_r(r; s)$ and the radial dependence of $Z[k; s, r, r']$, we decide whether to take the transverse average approach or keep the radial variable as parameters.

III. BUNCHED-BEAM THEORY

In reality, the electron beam's longitudinal distribution is not uniform, so we need to improve the theory to deal with a bunched beam. As mentioned above, with a nonuniform longitudinal distribution, a multifrequency theory is needed. This will be explored in this section.

A. One-dimensional formulae

For a 1D theory, the derivation up to Eq. (7) stays the same. For a Gaussian longitudinal distribution, we assume

$$\bar{f}_0(\mathbf{X}_0) = \frac{N}{2\pi\sigma_{\Delta\gamma}\sigma_z} \exp\left\{-\frac{\Delta\gamma^2}{2\sigma_{\Delta\gamma}^2} - \frac{z^2}{2\sigma_z^2}\right\}. \quad (31)$$

Here we use the same notation for $\bar{f}_0(\mathbf{X}_0)$ as in Eq. (9).

Completing the integral in Eq. (7), we obtain the evolution of the bunching factor as

$$b[k(s); s] = b_0[k(s); s] + \int_0^s d\tau \int \frac{dk(\tau)}{2\pi} K[k(\tau), k(s); \tau, s] b[k(\tau); \tau], \quad (32)$$

with the integral kernel

$$K[k(\tau), k(s); \tau, s] = ik(s)\mathcal{R}_{56}(\tau \rightarrow s) \frac{I(\tau)Z[k(\tau); \tau]}{I_A} \times \exp\left\{-\frac{[k(s) - k(\tau)]^2\sigma_z^2}{2} - \frac{[k(s)\mathcal{R}_{56}(s) - k(\tau)\mathcal{R}_{56}(\tau)]^2\sigma_{\Delta\gamma}^2}{2}\right\}. \quad (33)$$

The corresponding accumulated energy modulation spectrum is

$$\Delta\gamma[k(s); s] = -\int_0^s d\tau \int \frac{dk(\tau)}{2\pi} \frac{I_0 Z[k(\tau); \tau] b[k(\tau); \tau]}{I_A} \times \exp\left\{-\frac{[k(s) - k(\tau)]^2\sigma_z^2}{2} - \frac{[k(s)\mathcal{R}_{56}(s) - k(\tau)\mathcal{R}_{56}(\tau)]^2\sigma_{\Delta\gamma}^2}{2}\right\}. \quad (34)$$

It is now clear that the frequency is evolving along the beam line. Hence, it is intrinsically a multifrequency theory.

B. Radial dependence

We can further take the radial dependence into consideration. With the transverse averaging approach, we simply replace $Z[k(\tau); \tau]$ in Eqs. (32)–(34) by $\bar{Z}[k(\tau); \tau]$ defined in Eq. (22).

A similar derivation is easily obtained if we keep the radial variable as parameters. We omit detailed expressions here.

Once again, based on the detailed radial dependence, we have to choose either the transverse average approach, or to keep the radial variable as parameters.

IV. LSC IMPEDANCE WITH RADIAL DEPENDENCE

Having set up the framework in Secs. II and III, we work out some details in the next case where the impedance originates from the LSC effect. We first find the LSC impedance with a radial dependence.

A. Green function for a δ -ring

What we need is the Green function for a δ -ring. For a δ -ring at r' , the LSC impedance is derived to be

$$Z(k; r, r') = \frac{k}{\gamma} \left[\Theta(r' - r) \frac{2}{\gamma} K_0\left(\frac{kr'}{\gamma}\right) I_0\left(\frac{kr}{\gamma}\right) + \Theta(r - r') \frac{1}{kr'} K_0\left(\frac{kr}{\gamma}\right) \left\{ 2I_1\left(\frac{kr'}{\gamma}\right) + \frac{kr'}{\gamma} \left[I_0\left(\frac{kr'}{\gamma}\right) + I_2\left(\frac{kr'}{\gamma}\right) \right] \right\} \right], \quad (35)$$

where $\Theta(x) = 1$ for $x > 0$ and $\Theta(x) = 0$ for $x < 0$ is the Heaviside step function; K_0 , I_0 , I_1 , and I_2 are modified Bessel functions. We omit the derivation for Eq. (35) here.

For $f_r(r; s)$ being a parabolic distribution in r , one can get a closed form for the average impedance defined in Eq. (18). We omit the expressions for this closed form here.

V. COMPARISON

We now compare the results of our approaches to the results of other analytical approaches [6], and also to numerical simulations.

A. Comparison with other analytical approach

Compared to theories in Refs. [6, 7], our approach has the advantage of treating the real beam line, where the beam energy and the transverse beam size are varying, and the electron beam is bunched. Nevertheless, we make some comparison with Ref. [6] for a coasting beam with a constant energy and a constant transverse beam size. In their paper, they introduce a dimensionless parameter $q = k_m r_0 / \gamma_z$, with $k_m = 2\pi / \lambda_m$ where λ_m is the modulation wavelength; r_0 is the typical transverse beam size; and γ_z is the longitudinal Lorentz factor. Their theory reduces to 1D formulae when $q \rightarrow \infty$. Hence, we compare with the case of a small $q = 1$. In Figs. 1 and 2, we study the same examples as in Figs. 10, 11, 14, and 15 of Ref. [6]. The results are almost the same. For these studies, we show the radial dependence of the LSC-induced normalized longitudinal electric field \bar{E}_z . The initial longitudinal density modulation can be nonuniform transversely due to transverse granularity. For the case shown

in Fig. 1, the initial longitudinal density modulation is Gaussian transversely, *i.e.*, $\hat{a}_{1d} = e^{-\hat{r}^2/(2\sigma^2)}$ with $\sigma = 0.1$ and $\hat{r} \equiv r/r_0$. The induced \hat{E}_z is also nonuniform transversely. When the system evolves, the maximum \hat{E}_z can be shifted to $\hat{r} \neq 0$ quickly. For the case shown in Fig. 2, the initial longitudinal density modulation is transversely uniform. In this case, the transverse dependence of \hat{E}_z evolves much more slowly than in the case of Fig. 1. In these figures, \hat{z} is the normalized distance s along the beam line. We show \hat{E}_z at $\hat{z} = 0$, the beginning of the beam line, and at $\hat{z} = 10$, the end of the beam line for this calculation.

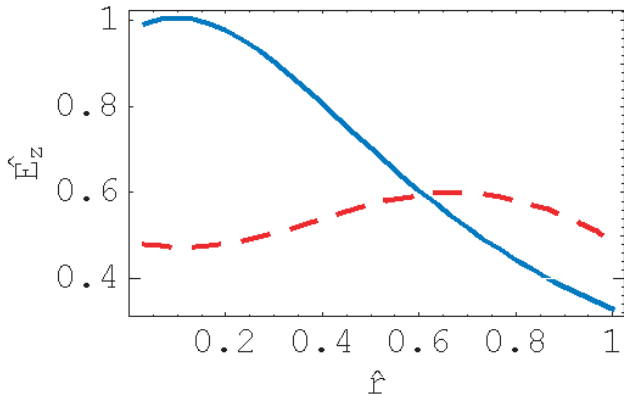


FIG. 1: Example of Figs. 10 and 11 of Ref. [6]. The initial density modulation is $\hat{a}_{1d} = e^{-\hat{r}^2/(2\sigma^2)}$ with $\sigma = 0.1$ and $q = 1$. The solid curve (blue) is at $\hat{z} = 0$, and the dashed curve (red) is at $\hat{z} = 10$. Notations have same meaning as in Ref. [6] and are defined in the text. The \hat{E}_z is normalized.

B. Comparison with PARMELA

Having compared our results to the results in Ref. [6] where they are applicable, we now deal with a realistic beam line. The example which we study is a beam with energy $E = 5.7$ MeV and peak current of 100 A. We study a 3 m long drift space, with betatron focusing, hence the rms transverse beam size σ_r varies from 0.5 mm to 3.7 mm. The initial density modulation is 5 % with wavelength of 0.5 mm. This yields $q \in [1.0, 7.2]$, taking $r_0 = \sqrt{3}\sigma_r$. We show in Fig. 3 the results of the four approaches developed in this paper. As a comparison, a PARMELA simulation [8] for a bunched beam, with rms bunch length $\sigma_z = 0.83$ mm, is also presented. Notice that in Fig. 3, $\Delta E = (\Delta\gamma)mc^2$ is the energy modulation.

VI. DISCUSSION

As shown in Fig. 3, the coasting-beam theory oversimplifies the LSC calculation in situations where the beam

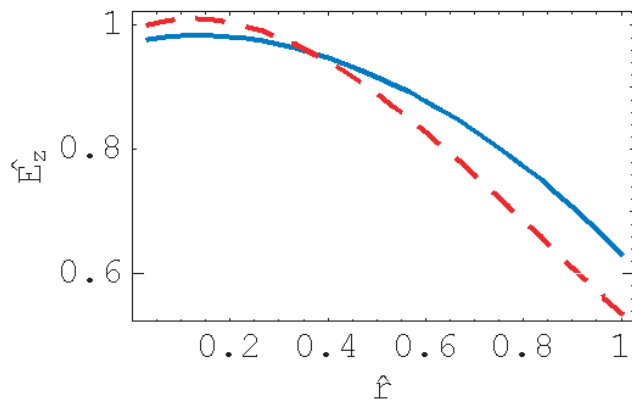


FIG. 2: Example of Figs. 14 and 15 of Ref. [6]. The initial density modulation is $\hat{a}_{1d} = 1$ with $q = 1$. The solid curve (blue) is at $\hat{z} = 0$, and the dashed curve (red) is at $\hat{z} = 10$. Notations have same meaning as in Ref. [6] and are defined in the text. The \hat{E}_z is normalized.

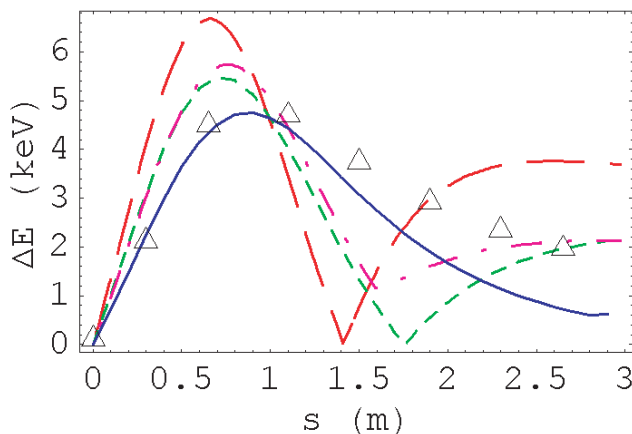


FIG. 3: Comparison of four different analytical approaches developed in this paper with PARMELA simulation. The long-dashed curve (red) is the 1D coasting-beam approach, the dashed curve (green) is the coasting beam with radial dependence approach, the dash-dotted curve (purple) is the 1D bunched-beam approach, the solid curve (blue) is the bunched-beam approach with radial dependence, and the “ Δ ” (black) is the PARMELA simulation.

is bunched. Even with radial dependence, the coasting-beam theory [6] can not capture some important features observed in simulations of a bunched beam. These profound features, showing mixing in the phase space, are crucial in answering questions of how fast the microbunching is damped or amplified along the beam line. This is particularly important in the front end of the accelerator, where the electron beam energy is low, and the longitudinal distribution function is not frozen. As shown, the analytical approach developed in this paper for a bunched beam with radially dependent impedance shows good agreement with the PARMELA numerical

simulations.

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