

Coherently Enhanced Radiation Friction in Laser-Plasma Collisions

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We reconsider the footprints of radiation friction in a head-on collision of a bunch of relativistic charged particles with a laser pulse by demonstrating that forward and backward radiation and radiation friction are coherently enhanced in a dense bunch. This opens an avenue to observe radiation friction in laser-matter interactions at much lower energies and laser intensities than previously considered possible. An estimate for the energy losses of the particles in the bunch over the collision due to coherent radiation friction in terms of laser and bunch parameters is derived and validated by comparing with the results of three dimensional particle-in-cell simulations.

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Introduction—A charged particle in an electromagnetic field accelerates, hence emits radiation, which eventually takes away its energy and momentum. In strong fields this might have a substantial impact on the particle trajectory. Besides slowing the particle down, this effect, conventionally called radiation friction (RF) or radiation reaction [1–3], can produce a highly nontrivial dynamics, including radiation trapping [4–7], the modification of the electron distribution in a bunch traversing a laser pulse [8–12] and the enhancement of longitudinal plasma waves in an underdense plasma [13,14]. RF can be also important in practical applications, such as electron and ion laser-plasma acceleration [8,9,15–21] or generation of strong magnetic fields in dense plasmas [22,23]. It also modifies the spectrum of radiation itself [10,24–26].

It has been conventionally believed that a substantial RF can arise either in a strong field or for high energy particles [24,27]. Strong electromagnetic fields can be generated by high power lasers [28]; therefore, RF attracts much attention in the context of laser–particle interactions; see, e.g., the reviews of Refs. [27,29–31]. More precisely, the energy of a particle is significantly altered by RF in a head-on laser-particle collision if [17,24]

$$\mathcal{R} = \mu a_0^2 \gamma \omega_L T \gtrsim 1, \quad \mu = \frac{4\pi r_e}{3 \lambda_L}, \quad (1)$$

where $a_0 = eE_0/m\omega_L c$ is the amplitude of the dimensionless field strength, m ; e and γ are the particle mass, charge, and Lorentz factor; ω_L , λ_L , and T are the pulse frequency, wavelength, and duration; and $r_e = e^2/mc^2 \approx 2.8 \times 10^{-13}$ cm is the classical electron radius. For an optical laser, $\mu \sim 10^{-8}$, implying that Eq. (1) is reached with $a_0 \gg 1$ and/or $\gamma \gg 1$.

Indeed, RF was observed experimentally [32–34] in collisions of ultrarelativistic electron beams (electron energy $\mathcal{E} \approx 0.25$ –2 GeV, $\gamma \approx 500$ –4000) with an intense laser pulse ($a_0 \approx 25$, $I \approx 1.3 \times 10^{21}$ W/cm², $T = 45$ fs, $\lambda_L = 800$ nm) ensuring $\mathcal{R} \sim 0.1$ –1. Such experiments are technically challenging because they require the overlap of ultrashort, tightly focused laser pulses with an ultrarelativistic electron bunch and control of interaction parameters that are hard to measure directly [35]. Alternatively, Eq. (1) can be met by increasing the interaction time T , e.g., with flying-focus pulses [36].

Here we suggest another strategy to observe RF in a collision of a laser pulse with a charged particle bunch based on a coherent enhancement of bunch radiation. Indeed, while the energy radiated incoherently by N particles scales as N , the energy of coherent radiation scales as N^2 [37–46]. Recently, we studied coherent radiation of an electron bunch driven by a strong counter-propagating pulsed electromagnetic plane wave [44], focusing mostly on high frequencies ($\omega \gg \omega_L$). Here we estimate the main contribution to RF, show that it comes from frequencies $\omega \sim \omega_L$ and derive its dependence on the laser and particle bunch parameters. We demonstrate both

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analytically and numerically that coherent radiation can provide a substantial RF even at moderate particle energies and laser intensities such that inequality in Eq. (1) is violated.

In this Letter we present the physical mechanism, scaling laws, and numerical validation of the coherently enhanced RF, and discuss prospects for its experimental observation. The companion paper [47] provides a detailed derivation of the analytical model and an extended analysis of its robustness.

Coherent radiation of a particle bunch—Consider radiation of a bunch of charged particles colliding with a laser pulse, the latter for simplicity represented by a pulsed circularly polarized plane electromagnetic wave. Figure 1(a) shows the distribution of radiated electric field after the collision; see End Matter for details of the numerical approach. The radiation is directed primarily along (forward scattering) and opposite (backward scattering) the laser propagation direction. The forward scattering leaves the incident laser frequency unchanged and corresponds to a diffraction on the bunch. As discussed below, such radiation from all the particles interferes constructively, resulting in a strong coherent enhancement.

It may seem counterintuitive that near-forward radiation, emitted against the particle motion, slows particles down. Although entirely classical, the effect can be clearly explained in terms of laser photons. Before scattering, they move longitudinally; afterward, their momenta acquire transverse components and angular spread; see Fig. 1(a). Since forward radiation has the incident laser frequency [see Figs. 1(b) and 1(c) and [47]], the photon energy is unchanged, so the gain in transverse momentum implies a reduction of the longitudinal component. The lost longitudinal momentum is partially transferred to the particles, slowing them down.

The backward radiation frequency ω is upshifted [44,48], we denote its wavelength $\lambda = 2\pi c/\omega$. For the particles in a bunch slice of thickness $d \lesssim \lambda$ oriented

perpendicular to the laser propagation direction, their backward radiation interferes constructively. If the bunch thickness $L > \lambda$, then the emission from adjacent slices of thickness $d \sim \lambda$ interferes destructively and the total backward radiation depends on the longitudinal shape of the bunch. For example, for a Gaussian shape high frequencies with wavelengths $\lambda < L$ are suppressed exponentially [40] and for a rectangular shape (uniform density for $0 < x < L$) as $\propto \omega^{-2}$ [44]; see also Ref. [47].

To calculate radiation of a particle bunch we start from an angular-frequency distribution of the emitted radiation by N particles [3],

$$\frac{d\mathcal{E}}{d\omega d\Omega} = \frac{\omega^2}{c} \left| \sum_{j=1}^N e_j \int \frac{dt}{2\pi} \mathbf{n} \times [\mathbf{n} \times \mathbf{v}_j] e^{i\omega(t - \frac{\mathbf{n} \cdot \mathbf{r}_j}{c})} \right|^2, \quad (2)$$

where $\mathbf{n}$ and ω are the direction and frequency of the emitted radiation and e_j , $\mathbf{r}_j(t)$, and $\mathbf{v}_j(t)$ are the charge, coordinates, and velocity of the j th particle of the bunch. It contains incoherent and coherent parts, and if all particles in the bunch have the same charges and initial velocities (these assumptions can actually be relaxed, see below), then the coherent and incoherent parts of the radiation spectrum can be represented in the form [40,44,47,49]

$$\frac{d\mathcal{E}_{coh}}{d\omega d\Omega} = \alpha N^2 \frac{d\mathcal{E}^{(1)}}{d\omega d\Omega}, \quad \frac{d\mathcal{E}_{incoh}}{d\omega d\Omega} = (1 - \alpha) N \frac{d\mathcal{E}^{(1)}}{d\omega d\Omega}, \quad (3)$$

where $d\mathcal{E}^{(1)}/d\omega d\Omega$ is the single-particle spectrum corresponding to Eq. (2) with $N = 1$ and $\alpha = |\langle e^{i\Phi} \rangle|^2$ is the modulus squared average of the particle initial phase factor. Coherent radiation dominates if $\alpha N \gg 1$ [44]; below, we assume this and neglect the incoherent part. The limits $\alpha = 0$ and $\alpha = 1$ correspond to fully incoherent and fully coherent radiation, respectively [44].

To calculate the coherence factor α , one has to specify the spatial shape of the bunch, i.e., the density distribu-

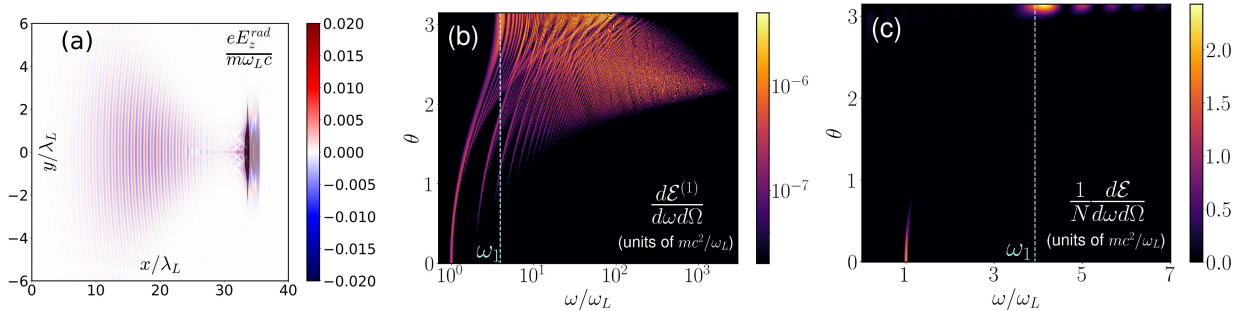


FIG. 1. Radiation in a head-on collision of electrons with a plane wave laser pulse. (a) the distributions of the z component of the radiated electric field in the (x, y) plane (in red and blue colors) and of the particle density (gray), the laser propagates along x axis to the left; (b) radiation spectrum of a single particle; (c) per-particle spectrum of a bunch of $N \approx 5 \times 10^6$ particles [see Eqs. (3) and (4)]; $\omega_1 = 2\pi c/\lambda_1$, is the fundamental harmonic of the backward radiation, corresponding wavelength λ_1 is defined below Eq. (9); $p_0 = 5mc$, $T \approx 50$ fs, $\lambda_L = 1 \mu\text{m}$.

tion $n(\mathbf{r})$ before the collision. For a Gaussian shape $n(x, \mathbf{r}_\perp) = n_0 \exp\{-[x^2/(L/2)^2] - (r_\perp^2/R^2)\}$, where R and L are the transverse and longitudinal sizes of the bunch,

$$\alpha = e^{-\frac{\omega^2}{8c^2}(L^2 \sin^2 \frac{\theta}{2} + 4R^2 \sin^2 \theta)} \quad (4)$$

selects narrow coherent peaks in the forward ($\theta = 0$) and, for sufficiently short bunches, backward ($\theta = \pi$) directions; for details see companion paper [47]. Here θ is the angle between the laser axis and the direction of the emitted radiation. Note that for a bunch smaller than the wavelength of the emitted radiation (both $\omega R/c \ll 1$ and $\omega L/c \ll 1$), we have $\alpha \approx 1$, meaning that all the particles in the bunch radiate coherently in all directions [41,44,50].

According to Eq. (3), the radiation spectrum of the whole bunch is proportional to the single particle radiation spectrum $d\mathcal{E}^{(1)}/d\omega d\Omega$, which can be obtained by numerical evaluation [51–53], as shown in Fig. 1(b) [54]. However, in forward and backward directions, where α peaks and therefore radiation is coherently enhanced, the analytical expression for the single particle spectrum can be evaluated approximately [44,48,58,59]; see companion paper [47]. In particular, in the forward direction the particle radiates at laser frequency, while in the backward direction the radiation frequency exceeds certain threshold $\omega_1 > \omega_L$. These features are clearly visible in Fig. 1(b) and are inherited by the bunch radiation; see Figs. 1(a) and 1(c).

Momentum transfer—When an ultrarelativistic particle collides head-on with an ultraintense laser pulse, it initially slows down, but then reaccelerates by the ponderomotive force. In the absence of RF the net change in longitudinal momentum $p_\parallel(t)$ (the component parallel to the laser pulse propagation) over the collision is zero, in accordance with the Lawson-Woodward theorem [60,61]. However, RF results in a net deceleration indicated by nonvanishing $\Delta p_\parallel \equiv p_\parallel(-\infty) - p_\parallel(\infty)$, which can be calculated as [17,24,47]

$$\Delta p_\parallel^{LL} \approx \sqrt{2\pi} \mathcal{R} p_0 \quad (5)$$

in accordance with Eq. (1). Here $p_0 = p_\parallel(-\infty)$ and the superscript LL indicates that the result is derived using the Landau-Lifshitz RF force.

Consider the average momentum loss $\langle \Delta p_\parallel \rangle = \Delta P_\parallel / N$ per particle in a bunch, where $\Delta P_\parallel$ is the total longitudinal momentum loss of the bunch. The particle momentum loss cannot be obtained by integrating the radiated energy Eq. (3) alone, since part of the emitted energy is drawn from the incident pulse. Instead we use conservation of the light-cone momentum $P_- = \mathcal{E}/c - P_\parallel$, for which the incident plane wave gives no contribution. Assuming that all the particles are ultrarelativistic during the entire course of the collision, so that $p_\parallel \approx -p_-/2$, this yields

$$\langle \Delta p_\parallel \rangle \approx \frac{1}{2cN} \int \frac{d\mathcal{E}_{coh}}{d\omega d\Omega} (1 - \cos \theta) d\omega d\Omega. \quad (6)$$

Since the angular distribution peaks at $\theta = 0$ and π , we split Eq. (6) into forward ($f, 0 < \theta < \pi/2$) and backward ($b, \pi/2 < \theta < \pi$) contributions. Both can be evaluated analytically when the radiation is narrow enough to use the single-particle spectra for exactly forward and backward scattering.

Details of the evaluation of these integrals are given in the companion paper [47]. As a result, the forward scattering contribution to the momentum loss reads,

$$\begin{aligned} \langle \Delta p_\parallel^f \rangle &\approx mc \frac{\pi^2}{8\sqrt{2}} \frac{n_0 a_0^2}{n_c \gamma_0^2} \omega_L T \mathcal{G}, \\ \mathcal{G} &= \frac{\sqrt{AB} (e^{A^2} - e^{B^2} + \sqrt{\pi} B [\operatorname{erfi}(B) - \operatorname{erfi}(A)])}{2\sqrt{2\pi} e^{B^2} \sqrt{B-A}}, \end{aligned} \quad (7)$$

where $A = \pi L^2 / [2\lambda_L \sqrt{2(16R^2 - L^2)}]$ and $B = 4\sqrt{2}\pi R^2 / (\lambda_L \sqrt{16R^2 - L^2})$, $\operatorname{erfi}(z)$ is the imaginary error function, $n_c = m\omega_L^2 / 4\pi e^2$ is the plasma critical density, γ_0 is the initial Lorentz factor of the particles before the collision, and we used $N = \pi^{3/2} n_0 R^2 L / 2$.

The dependence of RF on the bunch shape is encoded in the geometric factor $\mathcal{G}$. We discuss its properties and the corresponding features of the coherent radiation in the companion paper [47]. In particular, for a given $L \gtrsim \lambda_L$ the forward contribution Eq. (7) is maximized for

$$\frac{R^2}{L\lambda_L} \sim \text{const}, \quad (8)$$

where the value of the constant on the right-hand side is approximately 0.05. The corresponding maximum of the geometric factor is $\mathcal{G}^m \approx 0.03$.

Similarly, for the backward contribution we obtain [47]

$$\langle \Delta p_\parallel^b \rangle \approx mc \frac{\pi \omega_L T a_0^2 n_0}{32\sqrt{2} \gamma_0^2 n_c} e^{-\frac{\pi^2 L^2}{2\lambda_1^2}}, \quad (9)$$

where $\lambda_1 \approx [\lambda_L(1 + a_0^2)/4\gamma_0^2]$ is the wavelength of the fundamental harmonic of the backward radiation and we assumed that $2\pi L/\lambda_1 \gtrsim (\omega_L T)^{1/3} \gtrsim 1$. According to Eq. (9), the backward contribution to the momentum loss is strongly suppressed for $L \gtrsim \lambda_1$. In its derivation [47], we assumed that $R \gtrsim \lambda_1$. Therefore, as with the forward contribution Eqs. (7) and (9) may be inaccurate for a small bunch with both $R, L \ll \lambda_1$ [62]. For such small bunches an accurate result can be obtained by a direct numerical evaluation of the single particle spectrum $d\mathcal{E}^{(1)}/d\omega d\Omega$ in Eq. (6); see below. Recently, RF of such small bunches was

studied in Ref. [63] by numerical solution of the equations of motion for all particles in the bunch with interparticle interactions fully taken into account by means of Lienard-Wiechert potentials.

Numerical results—To test the theoretical predictions, we performed numerical simulations using particle-in-cell (PIC) code SMILEI [64]. The details of the numerical approach can be found in the End Matter.

We studied the dependence of the average longitudinal momentum loss on the bunch density n_0 , initial Lorentz factor γ_0 and the bunch sizes R and L . Figures 2(a)–2(c) compare Eqs. (7) and (9) with 3D PIC simulations, showing remarkable agreement. In Figs. 2(a) and 2(b) the simulation results start to deviate from the model predictions when $\langle \Delta p_{\parallel} \rangle / p_0 \gtrsim 0.01$. This is expected as the model relies on a single particle radiation spectrum computed along the particle trajectory in a sole laser field. When RF is large, it substantially modifies the trajectories and the coherent radiation yield should be corrected accordingly.

Figure 2(c) illustrates the dependence of the coherent RF on the bunch width R for two different values of thickness, $L = 2\lambda_L$ (thick) and $L = 0.2\lambda_L$ (thin). In the first case, the backward contribution is negligible and Eq. (7) [orange curve] accurately describes the effect for any R including the peak value $R \sim \sqrt{0.05L\lambda_L} \sim 0.3\lambda_L$. In the second case $L \ll \lambda_L$ the analytical estimate [turquoise curve] deviates for $R \lesssim \lambda_L$. This is natural since the assumption that radiation is sharply peaked around forward and backward directions is violated. However, by integrating Eq. (6) numerically with numerically evaluated single particle spectrum $d\mathcal{E}^{(1)}/d\omega d\Omega$, we found a very good agreement with PIC simulations even in this case, see the dark blue curve in Fig. 2(c). At the same time, for $R \gtrsim \lambda_L$ the sum of Eqs. (7) and (9) remains accurate.

The coherence of the emitted radiation can be limited by the bunch momentum spread [46,65]. To check the robustness of our results, we performed PIC simulations with a Gaussian momentum distribution $dN/d\mathbf{p} \sim e^{-(p_{\parallel}+p_0)^2/2\sigma_{\parallel}^2 - p_{\perp}^2/2\sigma_{\perp}^2}$, where $\sigma_{\parallel}$ and $\sigma_{\perp}$ are the initial

longitudinal and transverse momentum spreads. Figure 2(d) shows that RF is almost independent of the longitudinal momentum spread $\sigma_{\parallel}$, while the spread of the transverse momenta notably affects RF already at $\sigma_{\perp} \gtrsim 10^{-2}p_0$. The latter corresponds to the particle bunch divergence below 10 mrad, which has been achieved already in acceleration of high density bunches [66–69], see also Ref. [70].

The weak dependence on $\sigma_{\parallel}$ is expected, as for a thick bunch with $L \gtrsim \lambda_1$ the forward contribution to RF is dominant and at $\theta = 0$ the particle-dependent term in the phase in Eq. (2) vanishes [47]. Accordingly, the effect of momentum spread comes only from the spread in transverse components of the particle velocities in the preexponential factor of the integrand in Eq. (2).

Discussion—When a particle bunch collides with a laser pulse, the particles radiate coherently at low frequencies and incoherently at high frequencies [44]. For $a_0 \gg 1$ the incoherent radiation comes at $\omega \simeq a_0(a_0^2 + \gamma^2)\omega_L \gg \omega_L$ similarly to the single particle case [3,48,55]; see Fig. 1(b). In contrast, the coherent radiation is peaked at $\omega \sim \omega_L$ in the forward and at $\omega \sim \omega_1$ in the backward direction, respectively, and can be orders of magnitude stronger than the incoherent one, cf. Figs. 1(b) and 1(c).

Hence, for $a_0 \gg 1$ the coherent and incoherent contributions to RF are separated in the frequency domain. They can be compared using Eqs. (5) and (7). For the optimal bunch shape [see Eq. (8)], we have

$$\frac{\langle \Delta p_{\parallel} \rangle}{\Delta p_{\parallel}^{LL}} \sim \frac{n}{n_c} \frac{\mathcal{G}^m}{\mu \gamma_0^4}, \quad (10)$$

where μ is introduced in Eq. (1) and $\mathcal{G}^m \approx 0.03$. Since for an optical laser $\mu \sim 10^{-8}$, for a mildly relativistic bunch the coherent contribution should dominate already at bunch densities well below the critical density.

The dependence of the magnitude of RF on the bunch density and initial energy is illustrated in Fig. 3, where the

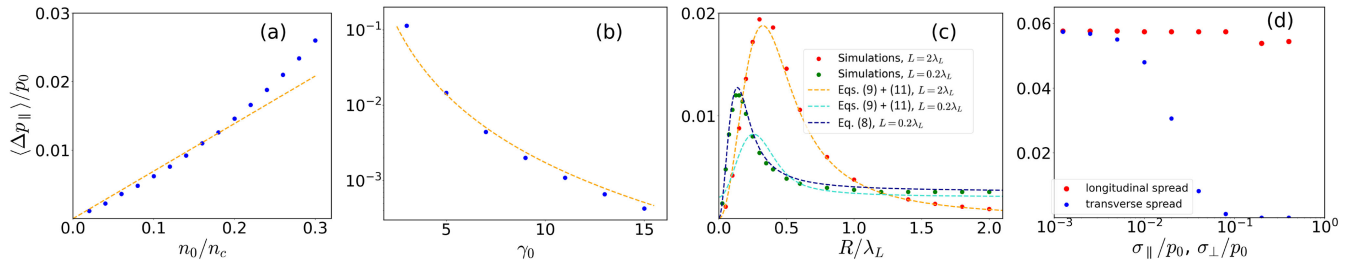


FIG. 2. Normalized particle deceleration $\langle \Delta p_{\parallel} \rangle / p_0$ over the collision vs (a) bunch density n_0 [for $\gamma_0 = 5, R = 0.5\lambda_L, L = 2\lambda_L$], (b) initial gamma-factor γ_0 [for $n_0 = 0.2n_c, R = 0.5\lambda_L, L = 2\lambda_L$], (c) bunch width R [for $n_0 = 0.2n_c, \gamma_0 = 5, L = 2\lambda_L, 0.2\lambda_L$], and (d) bunch momentum spread [for $n_0 = 0.5n_c, L = 2\lambda_L, R = 0.5\lambda_L$]. Blue, red, and green dots are extracted from 3D PIC simulations; dashed curves correspond to the sum of Eqs. (7) and (9) [orange/turquoise] and to the numerical integration of Eq. (6) with numerically calculated single particle spectrum [dark blue]; red and blue dots in (d) correspond to the longitudinal spread $\sigma_{\parallel}$ (for $\sigma_{\perp} = 0$) and the transverse spread $\sigma_{\perp}$ (for $\sigma_{\parallel} = 0$), respectively. Other parameters: $a_0 = 5, \omega_L T = 5\pi$.

ratio of the sum $\langle \Delta p_{\parallel} \rangle + \Delta p_{\parallel}^{LL}$ to the initial momentum p_0 is shown for the optimal bunch shape. At a lower laser intensity [$a_0 = 5$, $I \approx 6.6 \times 10^{19}$ W/cm²; see Fig. 3(a)] RF becomes substantial and potentially observable in two well separated regions at the top left and right corners of the plot. For lower bunch energy the major contribution to RF comes from coherent radiation at low frequencies and is proportional to the bunch density, while for a higher bunch energy it comes from incoherent radiation at high frequencies and is density independent.

At higher intensity [$a_0 = 25$, $I \approx 1.7 \times 10^{21}$ W/cm²; see Fig. 3(b)] RF remains substantial when the coherent and incoherent contributions are comparable (tilted green line). Here, however, one cannot expect Eqs. (5) and (7) to remain accurate, nor the total RF to be obtained as their simple sum, as their derivations assume that only one contribution dominates. Their joint regime is a separate interesting problem that we address for future studies.

Let us discuss the prospects for experimental observation of the coherently enhanced RF. The required laser intensities $I \gtrsim 10^{18}$ – 10^{20} W/cm² corresponding to $a_0 \sim 1$ – 10 and particle energies 10–100 MeV are now routinely achieved at the modern laser facilities [28]. However, production of relativistic particle bunches of the required high densities may be challenging. Obtaining electron bunches with densities of the order of 10^{18} – 10^{19} cm^{−3}, low emittance and energy spread $\delta\mathcal{E}/\mathcal{E} \sim 1$ –20% using laser/plasma wakefield accelerators were reported recently [66–68]. Maintaining a small energy spread is essential, because $\langle \Delta p_{\parallel} \rangle/p_0 \gtrsim \delta\mathcal{E}/\mathcal{E}$ is necessary for a straightforward observation of the effect. Even higher densities $n \sim 10^{21}$ cm^{−3} are anticipated at the FACET II facility in the near future [71,72] and in schemes based on electron acceleration by a helical laser [73–76]. For an optical laser ($\lambda_L \sim 1$ μm) these densities correspond to $(10^{-3} - 1)n_c$. For electron-positron bunches the state-of-the-art densities 10^{16} cm^{−3} [77] correspond to $10^{-5}n_c$ for an optical laser.

Note, however, that the ratio n/n_c could be also increased by lowering n_c , i.e., by using a larger wavelength

infrared laser instead of an optical one. For example, the critical plasma density corresponding to a 10 μm CO₂ laser is 100 times less than that for a 1 μm optical laser. The power of CO₂ lasers has already reached 10 TW level [78–80], thus making the required mildly relativistic regime $1 < a_0 < 10$ accessible.

It was initially unclear if radiation in experiments [32,33] was emitted by electrons in a classical or quantum regime. Only recently, after additional data was collected, the quantum regime of RF for that setup was confirmed [34]. In fact, observation of purely classical regime of RF in such experiments would be challenging, since the parameter in Eq. (1) ruling the significance of RF is proportional to the quantum parameter $\chi \approx 2\gamma a_0 \hbar \omega_L / mc^2$ [$\hbar$ is the Planck constant], which governs the importance of quantum effects [27,55]. For typical parameters of modern high power laser facilities [28] $a_0 \sim 10$ – 100 and $\omega_L T \sim 100$, the condition $\mathcal{R} \gtrsim 1$ implies $\chi \gtrsim 10^{-2}$ [see also Fig. 3], indicating that quantum effects are non-negligible. Observation of RF with mildly relativistic dense coherently radiating bunches instead of ultrarelativistic, rarefied ones requires much smaller values $\chi \sim 10^{-4}$ – 10^{-3} , thus giving access to a yet experimentally unexplored purely classical regime of RF.

To conclude, dense particle bunches can lose substantial energy and momentum in laser-bunch collisions via coherent radiation. Here we described analytically the dependence of the momentum loss by such a dense particle bunch colliding head-on with a laser pulse on the relevant laser and bunch parameters. This provides the first quantitative prediction for the scaling laws of coherent RF. Three dimensional PIC simulations confirm a high accuracy of our model. We emphasize that the coherently enhanced RF under consideration is not a correction to “ordinary” RF, which has been studied theoretically for decades and observed experimentally [32–34], but rather is a new mechanism of RF operating in a different parameter regime. The major contribution to coherent RF comes from forward and backward scattered radiation at low frequencies. Our findings demonstrate that observation of RF with dense particle bunches might be possible even at moderate particle energies and laser intensities. The coherent regime of radiation may also be relevant to the ongoing discussions of backreaction of the emitted radiation on the driving laser pulse [81,82], but we leave this for future studies.

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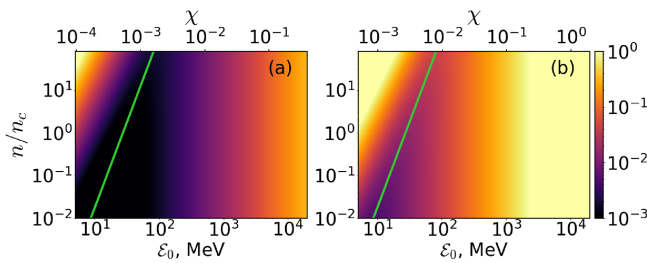


FIG. 3. Normalized particle deceleration ($\langle \Delta p_{\parallel} \rangle + \Delta p_{\parallel}^{LL} / p_0$) due to the coherent $\langle \Delta p_{\parallel} \rangle$ and incoherent $\Delta p_{\parallel}^{LL}$ contributions to RF for (a) $a_0 = 5$ [$I \approx 6.6 \times 10^{19}$ W/cm²] and (b) $a_0 = 25$ [$I \approx 1.7 \times 10^{21}$ W/cm²]. The green line corresponds to $\langle \Delta p_{\parallel} \rangle \sim \Delta p_{\parallel}^{LL}$; see Eq. (10). Other parameters: $\omega_L T = 15\pi$ [FWHM ≈ 50 fs], $\lambda_L = 1$ μm.

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Data availability—The data that support the findings of this article are not publicly available upon publication because it is not technically feasible and/or the cost of preparing, depositing, and hosting the data would be prohibitive within the terms of this research project. The data are available from the authors upon reasonable request.

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End Matter

Numerical approach—In a PIC simulation, at each time step the coordinates and momenta of macroparticles are updated according to the equations of motion in the instantaneous electromagnetic field. In turn, charge and current densities of macroparticles after smearing out across the cells serve as instant sources in the Maxwell equations governing the field evolution over the time step [83,84].

RF is conventionally accounted for by means of two major approaches [41]. The first one is just to add a RF force term to the RHS of the equation of motion for a macroparticle. This is considered reasonable when the quantum parameter

$$\chi \sim \frac{\gamma E}{E_{cr}}, \quad (\text{A1})$$

which measures the ratio of the electric field E in the particle rest frame to the critical field of quantum electrodynamics (QED) $E_{cr} = m^2 c^3 / e \hbar$ [85,86], is so small that the quantum effects of radiation can be neglected [27,55].

To account for quantum effects (including, e.g., straggling [87,88] and quenching [89]), instead of adding a classical RF force, one subdivides photons into soft and hard according to their energy, the latter also represented by macroparticles [90]. Their emission at each time step is then implemented as a random splitting off from charged macroparticles according to the probability distributions for photon emission in a locally constant crossed field approximation [41,55,91,92]. This way RF emerges as a quantum recoil of charged macroparticles ruled by energy-momentum conservation in a sequence of hard photon emissions. Soft photons are commonly ignored as their contribution to the ordinary (incoherent) RF can be neglected for strong fields ($a_0 \gtrsim 1$) and ultrarelativistic particles ($\gamma \gg 1$).

As they stand, both approaches account only for an incoherent contribution to RF. However, a coherent contribution is also partially captured in PIC simulations through the part of classical radiation of charged macroparticles that is resolved on the grid, i.e., radiation with wavelengths $\lambda \gtrsim d$, where d is the cell size [41]. This part of radiation correctly accounts for relative phases and hence for the coherence of radiating charged macroparticles, as in Eq. (2) [41,44]. With this RF emerges as a transfer of energy and momentum from macroparticles to their (coherent) radiation. This is seen in Fig. 4(a), where we compare simulation of the evolution of the average particle longitudinal momentum in the bunches of the same size but different densities with both the Landau-Lifshitz RF force and the quantum photon emission modules of the PIC code switched off. For dense particle bunches such contribution

to RF associated with coherent radiation at low frequencies $\omega \sim \omega_L$ can be essential and even dominant.

Note that qualitatively the evolution of the average momentum of a particle in a bunch is similar to the evolution of the momentum of a single particle subjected to Landau-Lifshitz RF force; cf. Figs. 4(a) and 4(b). However, it is worth emphasizing that due to coherent enhancement of radiation the particles in a bunch can lose a comparable amount of momentum for much lower values of energy and intensity.

Both analytical model and simulations presented above used the plane wave approximation of the laser field. This approximation should be valid, if $w \gg \lambda_L, R$ and $\lambda_R \gg cT$, where w is the laser waist and $\lambda_R = \pi w^2 / 2\lambda_L$ is the Rayleigh length. We performed an additional simulation with a focused laser pulse with $w = 4\lambda_L$ to demonstrate that the accuracy of the plane wave approximation in this case is better than 1%; see Fig. 4(a).

A dense relativistic electron bunch creates a strong self-field. This field increases the momentum spread in the bunch, but should not affect its total momentum, hence the average particle momentum. However, unless the solution of the relativistic Poisson equation at the initialization step of a simulation is perfectly accurate, a self-field can accelerate or decelerate the bunch as a whole and this is indeed what we observed in the simulations. To overcome the issue, we had to switch to simulations of neutral electron-positron bunches, for which such a numerical artifact is suppressed.

Fortunately, our approach works for bunches constituted of two species with the same masses and opposite charges, $|e_j| = e$. Indeed, in a circularly polarized field oppositely charged particles of equal masses orbit in the transverse plane with the same magnitude of velocity but in opposite directions. Therefore, for forward and backward scattering that make a major contribution to RF, the factor $e_j [\mathbf{n} \times [\mathbf{n} \times \mathbf{v}_j(t)]]$ in Eq. (2) is actually the same for all the particles in such a bunch, thus justifying Eq. (3) and the subsequent analysis. Hence, our theoretical predictions are equally valid for neutral electron-positron bunches, for which we compare them to the results of 3D numerical simulations.

In the simulations all the electrons and positrons constituting the bunch are moving from left to right with the same initial momenta [except the case of Fig. 2(d)], while the plane wave circularly polarized laser pulse is moving from right toward the bunch. The RF modules of the PIC code are switched off because the contribution of incoherent RF is negligible for the chosen simulation parameters and the deceleration is provided solely by the coherent radiation resolved on the grid. We extract from simulations a variation of the particle average momentum over the collision. The initial densities of electrons and positrons are given by the expression above Eq. (4) with the only replacement $n_0 \rightarrow n_0/2$ (so that n_0 now stands for the maximal value of the total particle density in the bunch). Each cell with plasma contains 5 macroparticles of each kind and the

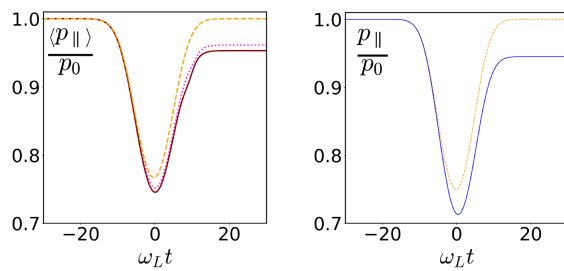


FIG. 4. (a) Evolution of the average longitudinal momentum of particles in a neutral electron-positron bunch colliding with a laser pulse obtained by numerical simulations for $a_0 = \gamma_0 = 5$, $T \approx 17$ fs, $R = 0.3\lambda_L$, $L = 1.5\lambda_L$, $\lambda_L = 1 \mu\text{m}$; solid dark red / dotted magenta lines: dense bunch $n_0 = 0.4n_c$, $N \approx 1.7 \times 10^8$, and plane wave/focused (waist $w = 4\lambda_L$) laser pulse; dashed orange line: rarefied $n_0 = 4 \times 10^{-5}n_c$, $N \approx 1.7 \times 10^4$ bunch and plane wave laser pulse; (b) evolution of the longitudinal momentum of a charged particle colliding with a laser pulse with (solid blue line) and without (dashed orange line) account for RF for $a_0 = \gamma_0 = 50$, $T \approx 17$ fs, $\lambda_L = 1 \mu\text{m}$.

simulation box size is $25\lambda_L \times 8\lambda_L \times 8\lambda_L$ for the simulations with plane wave pulses and $25\lambda_L \times 20\lambda_L \times 20\lambda_L$ for the simulation with the focused pulse. The spatial resolution is 64 cells per laser wavelength and the temporal resolution is

128 time steps per laser period. This resolution is sufficient to resolve the frequencies of the coherent radiation contributing to RF. The transverse and longitudinal boundary conditions are periodic and absorptive, respectively.