

Collective and surface charge effects in high-brightness ultrafast electron sources

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We present a photoemission model based on the Fowler-DuBridge equation that accounts for the dynamic modulation of the Schottky barrier during the emission of high-charge-density electron bunches. By incorporating the time-dependent effective electric field normal to the photocathode surface (including real-time emission-induced electric fields), our model captures the ultrafast transient modulation of the electron emission potential barrier that governs the instantaneous current density. This approach links single-electron photoemission effects with macroscopic collective charge effects, enabling accurate predictions of the time-dependent charge production for high-brightness photoinjectors. Furthermore, we validate the model by comparing it with the experimental results obtained with a high-gradient photoinjector equipped with ultraviolet pulse shaping capabilities. Our model highlights the importance of dynamic surface effects in shaping the initial phase space of electron bunches. The effects are particularly relevant for future lepton colliders where producing on-demand bunch charge with a tailored phase space is required, including the top-up injection scheme for CERN's Future Circular Collider (FCC).

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I. INTRODUCTION

Radiofrequency (rf) photoinjectors are routinely used to produce beams or bunches of electrons with the required characteristics for particle accelerators. Generally, the relevant parameters of performance are the spatial and temporal characteristics of the produced bunches, their charge and current density, as well as their momentum distribution (energy spread and beam brightness). Next-generation accelerator applications demand electron sources capable of producing ultrafast, precisely controlled bunches [1].

Photocathodes are routinely used to efficiently convert ultraviolet (UV) photons to electrons [2,3]. The electron bunch time-dependent current density characteristics are directly governed by the physics of the photoemission process, particularly for high-brightness electron beams [4,5]. Accurate modeling of these processes enables real-time control of the photoemission dynamics [6]. Accurate prediction of these emission dynamics requires models that

link the laser-driven photoemission process to the evolving near-surface electromagnetic fields.

In most simulation frameworks, emission is treated as an instantaneous process under static applied fields, since the characteristic timescales of field variation are typically much longer than those of the underlying emission mechanisms (tunneling, scattering, carrier transport, and radiative processes) [7]. These assumptions break down for ultrafast electron bunches, where fields and emission change on similar timescales. To fix this, recent developments using time-dependent quantum models (based on the Schrödinger equation or quantum distribution functions) can simulate how emission evolves dynamically, including delays, tunneling, and thermal effects [8,9]. However, these formulations do not include a self-consistent treatment of space charge or collective effects, which become significant when dealing with ultrafast UV pulses.

On the one hand, analytic frameworks describing electron transport through nonlinear potential barriers [10] extend classical thermal-field emission theory to more complex interface geometries, improving accuracy while remaining computationally efficient, though they also neglect time-dependent collective interactions. Numerical self-consistent methods, on the other hand, explicitly couple the evolving emission current with the electromagnetic fields near the photocathode. A self-consistent photoemission framework based on a three-dimensional Liénard-Wiechert particle-particle (LW-PP) algorithm has been demonstrated, in which

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the local cathode field dynamically modulates the Schottky barrier and emission yield [11]. While this approach achieves high physical fidelity, its $\mathcal{O}(N^2)$ computational scaling renders it too heavy for real-time applications where bunch production must be tailored on demand.

Particle-in-cell (PIC) simulation is a common approach for modeling the evolution of photoemitted electron bunches. An alternative self-consistent PIC-based approach introduces a transient positive charge layer within the photocathode to partially compensate the bunch space-charge field [12]. While this improves agreement with experiments, the method remains computationally intensive and self-consistent only at the macroscopic level, without resolving microscopic barrier dynamics. In a complementary line of work, a one-dimensional planar-gap model solved Poisson's equation self-consistently with a time-dependent emission current to analyze the onset of space-charge-limited flow [13]. This approach, while also macroscopic in nature, focuses on current limitation rather than the microscopic photoemission mechanism. However, since PIC methods require emission parameters as input and are computationally intensive, canonical photoemission models can serve as useful complementary tools, providing faster estimations of emissive behavior suitable for real-time machine studies and control.

A fully 3D, self-consistent point-to-point Coulomb model has also been developed to study ultrabright photoelectron beams [14]. This approach includes microscopic particle-particle interactions and image-charge effects to capture disorder-induced heating and emittance growth with high accuracy, but it is computationally intensive and focuses on postemission beam dynamics rather than the emission process itself.

Existing approaches either describe microscopic barrier physics under static conditions or resolve collective-field effects at high computational cost. In this paper, we present a compact time-dependent model that captures both the microscopic barrier modulation and macroscopic field evolution that are required to enable efficient, physically consistent simulations of transient emission.

The present model is intended to complement established particle-tracking approaches by providing a computationally efficient and physically transparent description of the time-dependent coupling between emission dynamics and Schottky barrier lowering.

II. FROM CANONICAL EMISSION LAWS TO SELF-CONSISTENT PHOTOEMISSION MODELS

Predicting electron bunch characteristics (and therefore the injector effective transfer function) is complicated by the dynamic superposition of near-cathode rf, image-charge, and space-charge fields. This evolving field environment transiently modulates the Schottky barrier and introduces nonlinear collective forces, leading to beam-quality degradation, most notably transverse emittance growth and phase-space distortions [15].

Within the broad landscape of postemission beam engineering strategies, one established approach is to generate an electron bunch with a uniformly filled ellipsoidal 3D charge distribution and exploit the linear self-fields to suppress emittance growth [16]. Another direction relies on premodulating the UV drive laser and allowing longitudinal space charge to wash out the initial density modulations, thereby producing a more linear longitudinal phase space [17]. In related injector-tuning paradigms, the space-charge-dominated beam is guided near laminar Brillouin flow in drifts and along the invariant envelope in accelerating sections to mitigate emittance growth. Recent work by Denham *et al.* has provided the first direct experimental validation of space charge saturation scaling laws in high-gradient rf photocathode guns, confirming the characteristic $Q_{\text{sat}} \propto (E_0 R)^{3/2}$ dependence, where E_0 is the accelerating electric field at the cathode and R is the effective transverse beam radius, and offering a practical framework for identifying the onset of extractable-charge saturation in modern injectors [18].

Complementary to these downstream control methods, one may instead address the problem at its origin by modeling the photoemission process itself, explicitly linking the UV laser spatiotemporal structure and photon energy to the time-dependent emission rate and initial phase-space distribution that ultimately shape injector performance.

Photoemission simulations typically follow two complementary tracks: (i) canonical/analytic emission equations that give fast, local estimates of current density and initial phase-space at the cathode, and (ii) PIC solvers that evolve emission and fields self-consistently on the device scale [19].

Canonical or analytic models include thermal photoemission (Richardson-Laue-Dushman) [20], field emission (Fowler-Nordheim or Murphy-Good) [21], and photoemission (Fowler-DuBridge) [22], as well as generalized thermal-field-photoemission formulations with semianalytic Wentzel-Kramers-Brillouin (WKB)/Gamow factors and barrier shape functions [23]. They are invaluable for fast, local estimates of steady or quasisteady emission under prescribed fields; however, by construction they neglect the time-dependent reshaping of the surface barrier induced by the evolving self-fields of the bunch (image and space charge) and nonlocal transport.

PIC-based emission solvers update particles and fields together, naturally capturing shielding, space-charge loading, geometric curvature of fields, and field-emission coupling under realistic device boundaries (with adaptive meshing where needed). PIC thus resolves macroscopic collective effects (including dynamic barrier reduction due to self-fields) and their feedback on emission, but still require a boundary emission law and, at practical resolutions, cannot resolve the microscopic barrier physics (image potential, tunneling through a time-tilted barrier) without external closure. At mesh sizes practical for device-scale PIC simulations, the code cannot resolve Å-nm barrier

physics—specifically the image potential and subfemtosecond tunneling through a time-varying reshaped barrier,

$$V(z, t) = \phi - e F(t)z - \frac{e^2}{16\pi\epsilon_0 z}. \quad (1)$$

Consequently, a physically based external closure must be supplied at the cathode boundary: an emission law $J[F(t), I(\mathbf{r}, t)]$ (e.g., Fowler-DuBridge with Schottky lowering and WKB/Gamow tunneling) that injects particles while the PIC self-consistently computes the device-scale fields feeding back into $F(t)$.

Taken together, neither track alone reproduces the specific effects most relevant here, namely: the transient Schottky modulation driven by the superposition of rf, image, and space-charge fields and its imprint on the instantaneous emission rate, tunneling probability, and initial phase space. As emphasized by Jensen: “*everything affects emission, and emission affects everything*” [19]: temperature, fields, and space charge are tightly coupled. A purely canonical treatment may miss device-level shielding and transient field evolution, whereas a pure PIC treatment without a microscopic emission kernel cannot reproduce barrier shape/tunneling specifics.

This motivates a hybrid, time-dependent micro-macro or multiscale treatment: we retain microscopic barrier physics while coupling it, at each time step, to the macroscopic field superposition at the cathode (nominal rf, surface-proximal image/dipole field from the nascent bunch, and finite-disk space-charge field). The result is a dynamic emission model that translates the UV driver’s spatiotemporal structure into physically consistent bunch distributions and observables (quantum efficiency, QE, current profile), while capturing collective feedback effects such as clamping and saturation that purely canonical or purely PIC approaches miss on their own.

III. THEORETICAL FRAMEWORK FOR DYNAMIC PHOTOEMISSION

We introduce a time-dependent extension to the Fowler-DuBridge (FD) photoemission model that explicitly incorporates the dynamic modulation of the Schottky barrier by evolving surface electric fields as shown in Fig. 1. By coupling instantaneous barrier lowering to the tunneling probability, our approach captures transient emission dynamics relevant for ultrafast, pulsed, and space-charge dominated regimes.

By bridging microscopic emission physics with macroscopic collective effects, the proposed model provides a predictive framework for the dynamic modulation of photoemission in high-brightness sources. Its capability to simulate the transient coupling between image-charge, space-charge, and external rf fields makes it particularly relevant for next-generation collider injectors, where

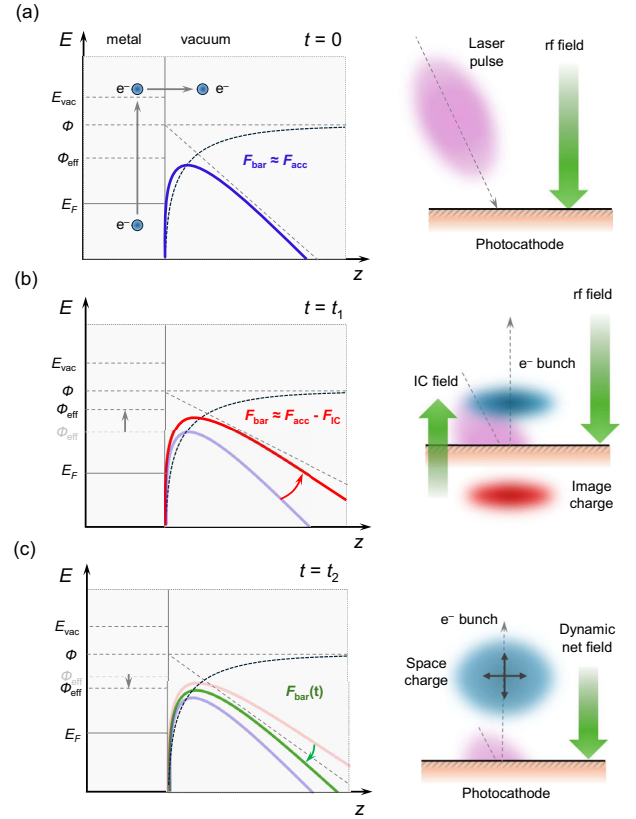


FIG. 1. Evolution of the effective Schottky barrier during photoemission from an rf-driven photocathode. (a) Initial condition ($t = 0$): the applied rf field lowers the effective barrier height ϕ_{eff} relative to the vacuum level E_{vac} , enabling electron emission from the photocathode surface. (b) Early emission ($t = t_1$): the emitted electron bunch generates an image-charge field that partially cancels the accelerating rf field, reducing the instantaneous barrier to $F_{bar} \approx F_{acc} - F_{IC}$. (c) Dynamic equilibrium ($t = t_2$): as the bunch propagates, space-charge and image-charge effects evolve, producing a net dynamic field that dictates the transient barrier height. The interplay of rf field, laser-driven emission, image-charge, and space-charge forces governs the time-dependent photocathode response.

producing on-demand, stable, and precisely tailored bunches is essential.

The coupled influence of the external accelerating field and the near-surface image-charge field determines the instantaneous shape and height of the emission barrier. As electrons are released and accumulate close to the surface, the local potential and effective work function evolve in time. Figure 1 schematically illustrates this dynamic balance: the accelerating field F_{acc} lowers the surface potential barrier via the Schottky effect, while the collective image-charge field F_{IC} partially counteracts this reduction. The resulting effective barrier field $F_{bar} \approx F_{acc} - F_{IC}$ evolves dynamically as emission proceeds, leading to a transient modification of the surface potential and of the effective work function ϕ_{eff} .

In the following, we formulate a time-dependent extension of the FD model that incorporates dynamic Schottky lowering and self-consistent collective feedback. The approach captures how the instantaneous superposition of rf, image-charge, and space-charge fields modulates the surface barrier and shapes the early bunch evolution. We begin by outlining the key physical mechanisms governing photoemission and their influence on projected emittance and then derive the field decomposition at the cathode surface used to separate barrier and transport fields. Subsequent subsections detail the finite-disk formulation of collective fields, the FD emission law with dynamic barrier lowering, and the computation of quantum efficiency under realistic laser and field conditions.

Photoemission from an rf-gun photocathode is governed by coupled material, optical, and electrodynamic effects: work function ϕ , image potential, surface morphology; photon energy $h\nu$, local fluence, and the spatiotemporal distribution $I(\mathbf{r}, t)$; the instantaneous rf surface field, which tilts the barrier and produces Schottky lowering; and space charge in vacuum and the photocathode's collective image response. We retain these ingredients with minimal structure. The laser provides $I(t)$ over an illuminated area; the external field is the nominal rf field at the cathode; collective effects enter through: collective image field that screens the barrier, and space-charge field that loads the early acceleration. This separates photoemission physics at the surface from beam dynamics in vacuum while avoiding double counting the collective effects. Here, we extend FD equation by letting the barrier follow the instantaneous surface field and by advancing the nascent bunch self-consistently.

The present model adopts a one-dimensional description of the emission process in which the electron potential well is assumed to have a fixed depth equal to the material work function. This approximation is appropriate for metallic photocathodes, where the high density of free electrons ensures rapid charge redistribution and efficient screening of any transient charge imbalance induced by emission.

As a result, the formation of a significant depletion region and associated modification of the internal potential is expected to be negligible. Under these conditions, the dominant time-dependent modification of the emission barrier arises from external fields, as captured by the Schottky barrier lowering included in the model. For materials with lower carrier density, such as semiconductors, internal electric fields and depletion effects may become significant and would require a self-consistent treatment of the potential inside the material. Such effects are not included in the present framework.

A. Finite-disk model and dynamic surface fields

Since photoemission is confined to the laser spot, the nascent bunch has a finite transverse extent. Nevertheless, most analytical injector treatments adopt the traditional

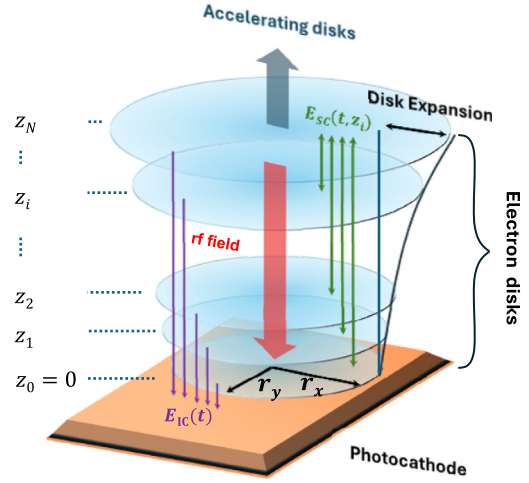


FIG. 2. Schematic of field decomposition during photoemission. The cathode at $z_0 = 0$ emits expanding electron sheets (radii r_x, r_y) driven by the rf field. Each sheet i feels a dynamic field $F_{\text{dyn}}(t, z_i) = F_{\text{rf}}(t) - E_{\text{SC}}(t, z_i)$ due to vacuum space charge, while the surface barrier field $F_{\text{bar}}(t) = F_{\text{rf}}(t) - E_{\text{IC}}(t)$ includes the collective image field. This separation avoids double counting and captures how image and space-charge effects shape early bunch evolution.

infinite-sheet (or “pancake”) approximation, in which the bunch is modeled as a transversely unbounded sheet of charge producing a constant field $E_z = \sigma/2\epsilon_0$ independent of distance from the photocathode [24–26]. While this pancake model captures the strong longitudinal compression at emission, it neglects the finite laser spot size and therefore overestimates the near-surface space-charge screening, artificially clamping the emitted current. Because the illuminated area is finite, the bunch field must decay with distance and depend on the spot radius, motivating the finite-disk model used here to describe both space-charge and image-charge fields self-consistently, as shown in Fig. 2.

For a uniformly charged disk of radius a and surface charge density $\sigma = q/(\pi a^2)$, the on-axis field a distance Δz away is

$$E_z^{\text{disk}}(\Delta z; a) = \frac{\sigma}{2\epsilon_0} \left(1 - \frac{\Delta z}{\sqrt{\Delta z^2 + a^2}} \right). \quad (2)$$

The finite-disk model therefore provides a consistent description of the surface field produced by the real and image sheets. To complete the formulation, we extend the same kernel to describe the vacuum space-charge forces acting on electrons once they leave the cathode. Using an identical disk representation for the emitted sheets ensures continuity between the surface barrier model and the in-vacuum dynamics, allowing the full longitudinal force law in Eq. (2) to be obtained.

During emission, we distinguish two complementary fields,

$$F_{\text{bar}}(t) \equiv F(z = 0^+, t) = F_{\text{rf}}(t) - E_{\text{IC}}(t), \quad (3)$$

$$F_{\text{dyn}}(t, z) \equiv F(z, t) = F_{\text{rf}}(t) - E_{\text{SC}}(t, z), \quad (4)$$

where F_{bar} is the surface (barrier) field entering Schottky lowering and tunneling, and F_{dyn} is the transport field in vacuum at the electron position z . $E_{\text{IC}}(t)$ is the collective image or dipole field produced at the surface by the already emitted charge that is still close to the cathode [27]. $E_{\text{SC}}(t, z)$ is the in-vacuum space-charge field evaluated at the particle location z . Using E_{IC} only in F_{bar} and E_{SC} only in F_{dyn} prevents overcounting.

Since the emission is represented as a sequence of thin disks, the cathode field is written as a sum over their individual contributions rather than as a continuous charge integral. Each disk has a closed-form finite-disk field, so superposition directly yields the total space-charge and image-charge fields.

Figure 2 summarizes the finite-disk geometry and its associated field decomposition. Each emitted sheet acts as a uniformly charged disk that expands in time, producing both a surface-proximal image field and a vacuum space-charge field. The figure illustrates how the barrier field $F_{\text{bar}}(t)$ and the transport field $F_{\text{dyn}}(t, z_i)$ are evaluated at different locations, providing a geometric interpretation of Eqs. (3) and (4) and clarifying the role of the finite-disk kernel in the near-cathode region.

For a uniformly charged disk of radius a and surface charge density σ placed at height $z = a$ above a grounded conducting plane, the method of images prescribes an image disk of identical radius and surface charge density $-\sigma$ located at $z = -a$ [28]. It is important to emphasize that the image disk contributes a field computed with the *same* finite-disk kernel as the real disk.

The barrier field is

$$F_{\text{bar}}(t) = F_{\text{rf}}(t) - E_{\text{IC}}(t),$$

$$E_{\text{IC}}(t) = \frac{1}{\epsilon_0 A} \sum_{j \in \mathcal{P}(t)} q_j \left(1 - \frac{z_j}{\sqrt{z_j^2 + a_j^2}} \right), \quad (5)$$

where $A = \pi r_x r_y$ denotes the UV laser spot area on the cathode, a_j is the instantaneous sheet radius (including transverse expansion), and z_j is its height. The fixed area $A = \pi r_x r_y$ reflects the UV laser footprint that sets the initial areal charge density, while the instantaneous expansion of each sheet enters through $a_j(t)$ in the finite-disk kernel. The image field is already accounted for: the real and image contributions combine to remove the usual $1/2$ factor in the finite-disk field, so the final expression requires no extra prefactor. The set $\mathcal{P}(t)$ contains only previously emitted sheets (no self-field).

While the finite-disk model determines the geometric strength of the image field, the actual field experienced

inside the cathode depends on the material response. In real metals, the image charge does not form an ideal mirror sheet at $z = 0$, but is instead distributed over a characteristic penetration depth. This behavior is described by an exponential attenuation factor $z_{\text{MAT}}(z)$, which modifies the surface field in Eq. (5).

$$E_{\text{IC}}(t) = \frac{1}{\epsilon_0 A} \sum_j q_j \left(1 - \frac{z_j}{\sqrt{z_j^2 + a_j^2}} \right) \times \exp \left[-\left(\frac{z_j}{z_{\text{MAT}}} \right) \right]. \quad (6)$$

Introducing this material term connects the idealized image-disk model to the realistic surface response. In the present formulation, this kernel is extended with an exponential fading factor $\exp[-(z/Z_{\text{mat}})]$, which accounts for the gradual loss of image-charge influence as electrons move a few micrometers away from the cathode surface. Physically, this term represents the finite spatial response of the metallic cathode: once electrons are sufficiently far from the surface, the induced mirror charge is no longer effective in sustaining the near-surface field. The characteristic decay length z_{MAT} defines the thickness of the region where the image-charge field remains significant. It represents an effective decay length associated with the persistence of image-charge-induced barrier modification following electron emission. Its value is expected to depend primarily on intrinsic material properties such as dielectric screening and surface electronic structure, as well as on factors such as surface quality or contamination. In the present model, z_{MAT} is treated as a material-dependent constant and is not varied with pulse duration.

The same finite-disk geometry is applied consistently to both the image-charge field E_{IC} and the space-charge field E_{SC} . The image field is evaluated at the cathode surface and includes the exponential fading factor to model the boundary response of the material, while the space-charge field is computed between emitted electron sheets within the vacuum region. This ensures a consistent near-field description of the collective fields acting on the emerging bunch.

The transverse laser profile is normalized using an elliptical Gaussian irradiance distribution characterized by effective radii r_x and r_y . Since the emission model is axisymmetric, the elliptical spot is mapped onto an equivalent circular sheet with the same illuminated area, such that the initial sheet radius is defined by

$$a_0 = \sqrt{r_x r_y}. \quad (7)$$

Thus, a_0 is an equal-area effective radius and should not be interpreted as a Gaussian rms or $1/e^2$ radius.

The transverse expansion of each sheet is treated ballistically from its initial radius a_0 , which is set by the laser spot size. In more complete electrodynamic

descriptions, the expansion is influenced by both the intrinsic thermal momentum spread and the transverse electric fields generated by the bunch itself [29,30]. However, transverse fields are not computed in the present model; the expansion is therefore governed solely by the transverse velocity associated with the mean transverse energy (MTE), $v_r = \sqrt{2\text{MTE}/m_e}$. With this approximation, the sheet radius evolves as

$$a_j(t) = \sqrt{a_0^2 + (v_r t)^2}. \quad (8)$$

The transverse velocity is initialized as $v_r = \sqrt{2\text{MTE}/m_e}$, with $\text{MTE} = 0.25$ eV taken from experimental measurements for metallic photocathodes [30–32]. As the sheet moves further from the cathode, its radius increases and its surface charge density $\sigma_j = q_j/(\pi a_j^2)$ decreases, reducing its contribution to both the image-charge and space-charge fields.

Although the image field acts only at the cathode surface, it is indirectly influenced by the evolving space-charge distribution above the surface. Space charge slows the longitudinal motion of later-emitted sheets, increasing their residence time near the cathode. This increases both their instantaneous height $z_j(t)$ and their transverse size $a_j(t)$, reducing their geometric contribution to $E_{IC}(t)$ through the factor $1 - z_j/\sqrt{z_j^2 + a_j^2}$.

As a result, the buildup of space charge moderates the image field experienced by later electrons and feeds back into the emission barrier. This coupling between longitudinal space-charge dynamics, transverse ballistic expansion, and image-charge screening provides a self-consistent link between the early beam evolution and the effective surface field governing photoemission.

As soon as electrons leave the cathode, the image field rapidly weakens and the bunch dynamics become dominated by the mutual repulsion of the expanding disks. Beyond the surface, the same finite-disk kernel can be used to evaluate the vacuum space-charge field produced by earlier sheets on those emitted later. This leads naturally to the following expression for the longitudinal field experienced by a sheet at position z_i . Particles advance with:

$$F_{\text{dyn}}(t, z_i) = F_{\text{rf}}(t) - E_{\text{SC}}(t, z_i),$$

$$E_{\text{SC}}(t, z_i) = \sum_{j: z_j < z_i} \frac{q_j}{\pi a_j^2 2\epsilon_0} \left(1 - \frac{z_i - z_j}{\sqrt{(z_i - z_j)^2 + a_j^2}} \right). \quad (9)$$

Equation (9) expresses the space-charge field acting on a sheet at position z_i is computed by summing contributions from sheets located closer to the cathode, i.e., with $z_j < z_i$.

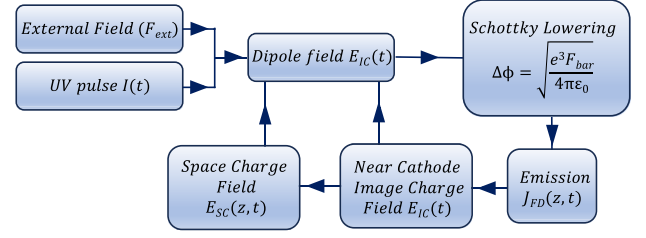


FIG. 3. Sequence of space-time-dependent steps involved in photoemission in this model.

Each term accounts for the finite transverse extent of the bunch through the disk radius a_j , which reduces the near-surface field compared with the infinite-sheet approximation. The resulting $E_{\text{SC}}(t, z_i)$ opposes the accelerating field, slowing newly emitted electrons. Together with the image-charge term in Eq. (6), this formulation provides a self-consistent description of how collective effects shape both the emission barrier and the early beam evolution.

This framework establishes the foundation for the coupled treatment of image and space-charge effects at the cathode surface, which is schematically summarized in Fig. 3. We note that the flowchart shown does not depict the microscopic emission physics itself, but rather the numerical structure by which our model translates the underlying physical couplings into solvable components. It serves as a conceptual map of how the external, image-charge, and space-charge fields are treated and updated within the simulation framework, rather than as a literal sequence of physical processes.

B. Modified time-dependent Fowler-DuBridge formalism

The microscopic image appears only in the single-electron barrier. The effective work function is lowered by the Schottky term:

$$\Delta\phi_S(F) = \sqrt{\frac{e^3 F}{4\pi\epsilon_0}}, \quad (10)$$

with $\phi_{\text{eff}}(t) = \phi - \Delta\phi_S(F_{\text{bar}}(t))$.

For metallic cathodes such as copper, photoemission is governed by a threshold behavior set by the work function, and the quantum efficiency exhibits a relatively weak and smooth dependence on wavelength above threshold. In contrast, semiconductor photocathodes show a much stronger wavelength dependence due to their band structure and electron affinity, which can lead to significantly enhanced quantum efficiency near optimal excitation wavelengths. The present model focuses on a metallic cathode and does not include these semiconductor-specific effects.

The FD current density can be calculated as

$$J_{\text{FD}}[F_{\text{bar}}(t)] = A_{\text{FD}}[\hbar\nu - \phi_{\text{eff}}(t)]_+^2 \exp\left(-\frac{B}{F_{\text{bar}}(t)}\right) I(t). \quad (11)$$

Here, A_{FD} is a material-dependent proportionality constant that scales the photoemission yield and is typically calibrated from low-fluence experimental data. The coefficient B controls the exponential field dependence of emission, corresponding to the barrier-transmission factor in the FD formalism. The term $I(t)$ is the laser intensity. If $\hbar\nu < \phi_{\text{eff}}(t)$, we apply a WKB transmission through the rf-tilted, image-corrected barrier,

$$V(z; t) = \phi - eF_{\text{bar}}(t)z - \frac{e^2}{16\pi\epsilon_0 z}, \quad (12)$$

$$T(t) = \exp\left[-\frac{2}{\hbar} \int_{z_1}^{z_2} \sqrt{2m_e V(z, t)} dz\right]$$

and the emitted current density is $J(t) = J_{\text{FD}}(t)T(t)$. The explicit image potential in (12) is single electron; the macroscopic image/SC fields are handled only via (3) and (4).

For a laser pulse of energy E_L , full width at half maximum (FWHM) duration τ , and elliptical $1/e^2$ radii (r_x, r_y) , the temporal irradiance profile is modeled as a Gaussian.

Assuming a Gaussian temporal profile and Gaussian transverse irradiance distribution, the laser intensity is written as [33]:

$$I(x, y, t) = I_0 \exp\left(-\frac{t^2}{2\sigma_t^2}\right) \exp\left(-\frac{2x^2}{r_x^2}\right) \exp\left(-\frac{2y^2}{r_y^2}\right), \quad (13)$$

where $\sigma_t = \tau/(2\sqrt{2\ln 2})$ and τ is the pulse duration (FWHM). Integrating over space and time yields the pulse energy

$$E_L = \iiint I(x, y, t) dx dy dt = I_0 \sigma_t A \sqrt{\frac{\pi}{2}}, \quad (14)$$

with $A = \pi r_x r_y$ denoting the effective elliptical spot area. The peak irradiance is therefore given by

$$I_0 = \sqrt{\frac{\ln 2}{\pi}} \frac{4E_L}{\tau A}. \quad (15)$$

The quantum efficiency (QE), defined as the ratio of emitted electrons from the rf gun to incident photons, is given by

$$\text{QE} = \frac{1}{N_\gamma} \int J(t) A dt, \quad N_\gamma = \frac{E_L}{\hbar\nu}, \quad (16)$$

where N_γ is the total number of photons in the laser pulse, and $J(t)$ is the time-dependent emission current density computed from Eqs. (11) and (12) using the instantaneous barrier field $F_{\text{bar}}(t)$.

Intuitively, the buildup of near-surface charge increases the collective image field $E_{\text{IC}}(t)$, which reduces the effective barrier field $F_{\text{bar}}(t)$ and weakens the Schottky lowering $\Delta\phi_S$. Simultaneously, the space-charge field E_{SC} develops in the vacuum region, opposing the extraction field and potentially recapturing slower electrons. Together, these feedback effects reduce the overall emission efficiency at high fluence, leading to a quantum efficiency roll-off, or clamping. Short, intense pulses enhance this effect through strong instantaneous E_{IC} , whereas larger laser spots mitigate it by distributing charge over a wider area, thereby reducing both E_{IC} and E_{SC} .

For the simulations, we discretize the laser pulse and evolve macro-sheets: at each time step, compute F_{bar} and F_{dyn} from (3) to (4) using (5)–(9); evaluate J_{FD} at $F_{\text{bar}}(t)$ and apply $T(t)$; convert to $dq = J A dt$ and append a sheet at $z = 0^+$ with initial normal velocity set by the MTE and transverse radius advanced by thermal expansion; advance all sheets with $a = eF_{\text{dyn}}/m_e$. A single calibration of A_{FD} is performed at weak fluence and fixed field; thereafter, the model predicts $Q(E_L)$ and QY vs fluence across spots and pulse lengths.

The barrier includes the explicit single-electron image term and an rf-tilted vacuum level; in-solid transport is absorbed into A_{FD} . The rf field is taken constant over the ps-scale pulse (generalization to $F_{\text{rf}}(t)$ is straightforward). Early transport is computed on-axis with finite-disk geometry for the source sheets; transverse phase space may be added in future work.

The proposed framework is particularly advantageous in regimes requiring rapid evaluation and in situations where analytical insight into ultrafast emission dynamics is essential.

The overall structure of the numerical implementation is summarized in Appendix B.

C. Implementation and computational performance

The model was implemented in Python and benchmarked on a standard workstation with 16 GB RAM. A full charge-fluence curve required approximately 0.13 s on average, corresponding to about 14 ms per simulation.

The numerical resolution was assessed using $N_t = 42$ temporal slices across the laser pulse. Increasing the temporal resolution by 50% ($N_t = 63$) changed the final simulated bunch charge by less than 1.2%. A summary of all parameters used in the model is provided in Table I.

IV. EXPERIMENTS

This section describes the experimental setup used to generate, transport, and characterize electron bunches in the S-band rf photogun. The system integrates a UV-driven photocathode source, synchronized laser and rf subsystems, and a complete suite of beam diagnostics. Particular attention is given to the laser system, including its rf synchronization and ultraviolet frequency conversion, as well as the optical transport and conditioning line that delivers the beam to the photocathode. The electron beam generation, acceleration, and measurement stages are then outlined, providing the experimental basis for the subsequent comparison with the theoretical model. The developed model was applied to simulate the photoemission behavior of the copper photocathode for the AWAKE run 2 c experiment witness bunch injector at CERN [34].

A. Laser and optical systems

A femtosecond ytterbium-doped potassium gadolinium tungstate (Yb:KGW) laser (PHAROS, Light Conversion) is used both for electron bunch generation. The laser delivers 1030 nm pulses with a maximum energy of 2.1 mJ and a minimum pulse duration of 190 fs FWHM. Although it can operate up to 3 kHz, the repetition rate was limited to 10 Hz for these experiments. A general layout of the laser and optical arrangement is shown in Fig. 4.

The synchronization between the laser oscillator and the rf system is achieved using a *PhaseLock* module (TEM Messtechnik). The system operates with two input signals: a 1.5 GHz reference derived from the master rf source and a 75 MHz signal from a high-bandwidth photodiode monitoring the laser oscillator output, which itself operates at 75 MHz. The 1.5 GHz signal, corresponding to the 20th harmonic of the laser repetition rate, is filtered and amplified before being compared with the rf reference.

A frequency counter decimates the 1.5 GHz reference to generate the 75 MHz signal used to perform rf bucket control. In this configuration, the laser can be tightly phase locked to the accelerator rf with excellent temporal stability, while maintaining the capability to introduce controlled phase shifts and scan large delay ranges as required for the experiments. The output of the phase-locked loops (PLLs) is used to control two piezoelectric actuators (PZT₁ and PZT₂), the one with larger displacement is used for slow drifts (low frequency loop), while the second is used for fast and small fluctuations of up to 1 kHz. The PLL compares the phase of the photodiode signal derived from the light conversion PHAROS pulse train with external 75 MHz and 1.5 GHz rf references, generating an error signal proportional to their phase difference. This signal is filtered and amplified to drive the PZTs actuators mounted on a cavity mirror, thereby adjusting the optical path length and stabilizing the laser repetition rate to the reference.

The laser pulse duration is continuously tunable between 190 fs and 15 ps via a motorized grating compressor, allowing for chirped-pulse operation. The laser output is frequency-quadrupled to 257 nm through two successive BBO crystals. The first, a 2-mm-thick crystal, converts the infrared beam to 515 nm with an efficiency of 55%, delivering 1.2 mJ pulses with 0.04% rms energy fluctuation. The second stage employs a 500 μ m BBO crystal to generate the fourth harmonic at 257 nm with 35% conversion efficiency, yielding 480 μ J of UV energy and 0.09% rms stability. The UV pulse duration is presently controlled by chirping the infrared pulse within the laser system, which reduces the overall conversion efficiency but enables tunability between 265 fs and 1.5 ps (FWHM) while maintaining pulse energies above 35 μ J. Pulse durations at the injection point were characterized using an APE two-photon autocorrelator [35].

The UV laser pulse is directed onto the copper cathode to generate electron bunches. The laser beam is transported from the laser room, located above the accelerator bunker, to the photocathode through a sequence of mirrors and periscopes spanning an overall optical path length of approximately 15 m from the point of UV conversion to the cathode surface. The beam is directed onto an optical table in the vicinity of the rf gun, where further conditioning takes place.

In this area, a motorized telescope allows on-demand adjustment of the beam magnification and resulting UV laser spot size on the photocathode from approximately 100 μ m in diameter ($1/e^2$) to 5 mm continuously. Since the system is not strictly image relayed onto the cathode, the transverse beam profile can vary between Gaussian-like and near flat-top shapes depending on the telescope configuration. For the present experiments, the telescope was tuned to produce Gaussian-like spots, with elliptical profiles obtained by adjusting the lens positions and tilt. Ellipticity in the UV beam can originate from phase-matching effects in the BBO crystal, leading to spatial walk-off, and from astigmatism introduced by off-axis optical elements in the beamline.

The transverse laser profiles were monitored using a diagnostic line replicating the virtual photocathode plane. Here, a 50- μ m-thick Nd:YAG scintillator converts the UV light to the visible range, which is then re-imaged onto a CMOS camera (Basler) via an achromatic 1:1 telescope. These images were used to characterize the spatial intensity distribution and to provide input parameters for numerical simulations. In the same diagnostic area, a calibrated energy meter was installed to measure the shot-to-shot pulse energy delivered to the photocathode. This combination of optical controls and diagnostics enables flexible manipulation of the three-dimensional photon distribution impinging on the cathode across a wide parameter range.

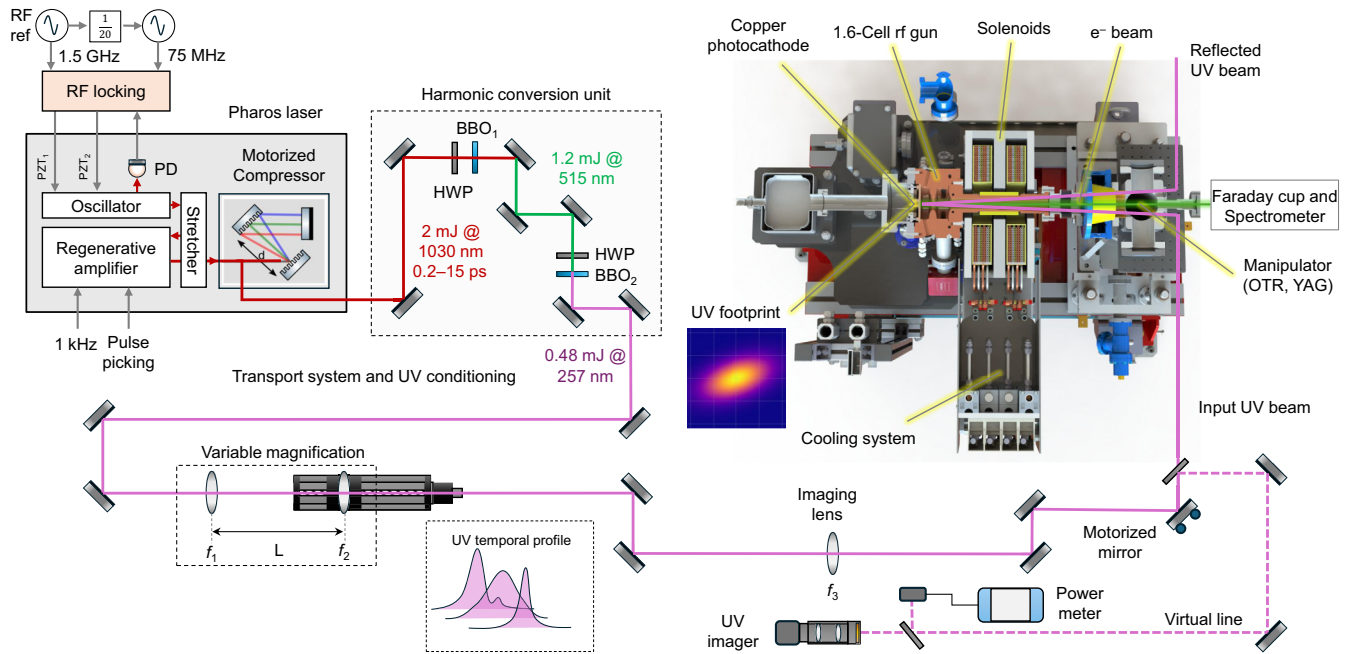


FIG. 4. Schematic of the laser and optical transport system driving the rf photogun. The Yb:KGW PHAROS laser is phase locked to the 1.5 GHz rf reference. Pulse duration is tuned between 190 fs and 15 ps via a motorized grating compressor. The 1030 nm beam is frequency-quadrupled to 257 nm through two BBO stages and transported over a 15 m optical path to the photocathode. The conditioning system includes a motorized telescope (lenses with focal lengths f_1 and f_2 at a variable distance L) to control the spot size, an imaging lens (focal length f_3), and a steering mirror for active control of spot alignment. A pick-off mirror is used for reflecting a small portion of the laser beam, which is directed to a diagnostic line with a Yb:YAG scintillator and CMOS camera records the transverse UV profile at a virtual cathode plane, while a calibrated energy meter monitors the pulse energy delivered to the photocathode. PZT: piezoelectric actuator, PD: photodiode, HWP: half-wave-plate. For electron beam production and measurement: a temporally and spatially shaped UV beam impinges a copper photocathode. The resulting electron beam is then refocused and characterized using a solenoid and screens in a manipulator arm. The charge and spectrum of the electrons are characterized using a Faraday cup and a magnet spectrometer.

B. Electron beam production and characterization

The electrons are produced in an S-band 1.6-cell rf photogun operating at 2.9985 GHz, fabricated in INFN Frascati based on the BNL/UCLA/SLAC design and using brazing-free technology [36]. The gun operates at a peak accelerating gradient of 106–120 MV/m. It delivers electron beam energies up to 6 MeV with a typical relative energy spread of 0.2%–2%. The photogun is equipped with a copper cathode, which minimizes thermal emittance compared to alternative semiconductor photocathodes. Two optical injection ports provide access for the drive laser: one introduces the beam at a 70° incidence angle through side viewports near the cathode, while the second delivers light at a shallow 3° angle using an off-axis mirror positioned roughly 0.5 m from the cathode outside the vacuum chamber, which was used in our experiments to position the UV beam on the photocathode center.

Single-electron bunches are generated at repetition rates between 0.83 and 10 Hz, with a nominal extracted charge of 400 pC, although up to 800 pC has been achieved for the shortest UV laser pulse duration of 190 fs. The copper photocathode typically exhibits a quantum efficiency of

around 10^{-4} to 10^{-5} , and measurements indicate an intrinsic emittance of 0.7 mm mrad/mm. The stability of the electron bunch charge was directly linked to the stability of the UV pulse at the photocathode. Since the UV beam was transported in air, cumulative nonlinear effects degrade its energy stability from 0.09% rms at the quadrupler output to approximately 1.04% at the rf gun input.

Immediately downstream of the gun, a four-coil solenoid assembly about 20 cm long provides transverse focusing. This element plays a central role in optimizing the beam brightness by producing a double emittance minimum along the transport line. Following the solenoid, a compact diagnostic section allows characterization of the beam in terms of its transverse profile, energy, charge, and emittance.

After acceleration, the solenoid focuses the beam toward a viewport chamber located 0.57 m from the cathode, where it is imaged using a YAG:Ce screen and a CMOS camera. This imaging system enables precise measurement of beam size and position as functions of key tuning parameters, including rf phase, laser spot size and position, and solenoid strength. Downstream, a magnetic dipole

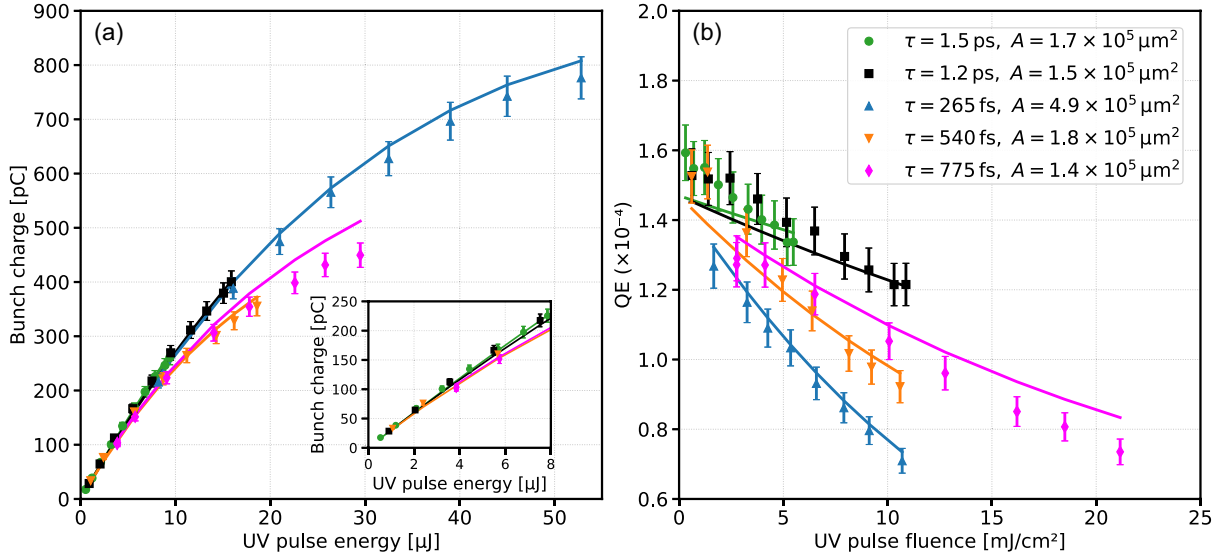


FIG. 5. Simulated (solid lines) and experimental (markers) of bunch charge vs UV pulse energy (a) and quantum efficiency vs fluence for deep-UV pulses (b) with different illumination parameters at the cathode surface.

spectrometer measures the electron beam momentum, and a Faraday cup records the bunch charge, with a stability better than 5%.

C. Results

The simulation reproduced the experimental laser and field parameters, including the measured pulse energy range (5–30 μ J), extraction field ($F_{\text{ext}} = 120$ MV/m), and spot geometries. The emission process was modeled through the time-dependent Fowler-DuBridge formalism, incorporating the collective image field $E_{\text{IC}}(t)$ and the evolving space-charge field $E_{\text{SC}}(t, z)$, as described in Sec. III. The photoemission yield and charge evolution were computed by integrating the instantaneous current density over the UV temporal envelope and cathode area.

Figure 5 compares the simulated and measured bunch charge as a function of UV pulse energy for several pulse durations. The model reproduces the characteristic sub-linear dependence of charge on fluence, consistent with experimental observations, capturing the onset of space-charge-induced saturation at high laser energies. Shorter pulses (265 fs) exhibit higher peak emission rates and earlier charge clamping, while longer pulses produce proportionally smaller space-charge fields, resulting in a more linear charge scaling. The corresponding simulated quantum efficiency curves showed good agreement with experimental data, validating the collective-field coupling in the model. Beyond the nominal saturation charge, additional emission can occur from the lower-intensity tails of a nonuniform laser spot, even when the central cathode region is fully saturated. This “tail-emission” effect leads to a slow, logarithmic increase of extracted charge with laser energy, which can be captured by tail-emission models [18,37].

To further contextualize the level of agreement with experimental observations, we note that established particle-tracking simulations are, in many cases, capable of reproducing similar trends across a broad range of operating conditions. In such approaches, emission suppression arising from the interplay of image-charge and space-charge forces is captured implicitly through particle reflection at the cathode surface, which can lead to an effective reduction in the extracted charge that mimics a dynamical quantum efficiency.

The simulation provides direct access to the transient barrier field $F_{\text{bar}}(t)$, effective work function $\phi_{\text{eff}}(t)$, and emission current density $J(t)$ for different laser pulse durations. For the 265 fs pulse ($F_{\text{ext}} = 120$ MV/m, spot $335 \times 469 \mu\text{m}^2$), the high peak fluence produces a rapid buildup of the collective image field $E_{\text{IC}}(t)$, which reaches approximately 20 MV/m within 100 fs. This transient screening of the applied field suppresses mid-pulse emission and leads to a strong asymmetry in $J(t)$, consistent with the observed saturation of the bunch charge around 700–800 pC.

In contrast, for the 1.5 ps pulse ($139 \times 393 \mu\text{m}^2$), the slower temporal envelope allows the image field to rise more gradually, reaching only 8–10 MV/m near the pulse peak. The resulting emission follows the laser profile more closely, with weaker clamping and a nearly linear scaling of the bunch charge up to ≈ 250 pC. Together, these two regimes illustrate how the interplay between the external field and the dynamically evolving image charge governs the onset and strength of emission saturation in ultrafast photoemission.

V. DISCUSSION

To illustrate the time-dependent behavior captured by the model, detailed simulations were performed for the two

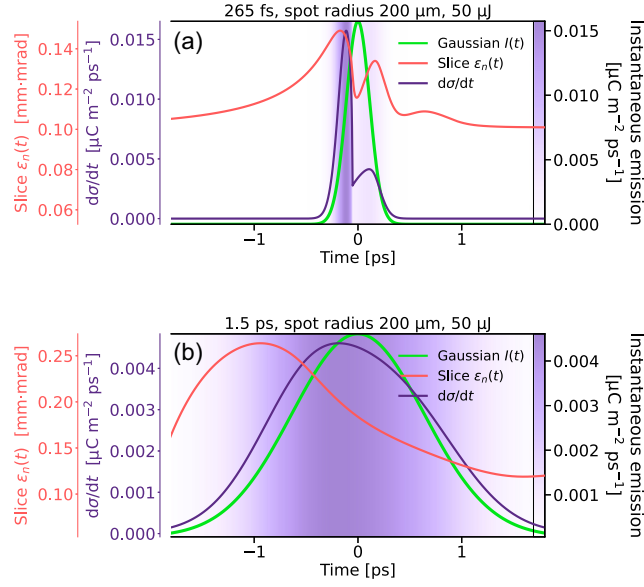


FIG. 6. Time-resolved emitted charge density and instantaneous (slice) normalized emittance $\epsilon_n(t)$ for two pulse durations. Background heat map: instantaneous emission $d\sigma/dt$ [$\mu\text{C m}^{-2}\text{ps}^{-1}$]; green curve: normalized laser profile $I(t)$; red: slice emittance $\epsilon_n(t)$ [mm mrad]; purple: emission rate (colorbar: absolute scale). Top (265 fs): rapid image-charge suppression produces early and late emittance peaks. Bottom (1.5 ps): extended emission shifts the emittance maximum to the leading edge, followed by monotonic decay.

extreme pulse durations in the dataset—the shortest (265 fs) and longest (1.5 ps) laser pulses. For a consistent comparison, both cases were simulated with the same pulse energy (50 μJ) and identical effective spot size ($200 \times 200 \mu\text{m}^2$), so that differences in emission arise solely from the temporal envelope of the laser. This setup isolates the influence of pulse duration on the buildup of the image-charge field, barrier modulation, and the resulting emission rate and emittance evolution. The following figures present the simulated instantaneous and cumulative quantities—including the image-charge field, quantum efficiency, emitted charge density, and slice emittance—providing direct insight into the self-consistent feedback mechanisms that govern ultrafast photoemission.

Figure 6 presents the time-resolved evolution of the emitted charge density and slice normalized emittance for the two pulse durations. The background heat map encodes the instantaneous emitted charge density $d\sigma/dt$ [$\mu\text{C m}^{-2}\text{ps}^{-1}$] at the cathode. The green curve is the Gaussian laser profile $I(t)$ (normalized). The red curve is the slice normalized emittance $\epsilon_n(t)$ [mm mrad] computed from the set of emitted sheets present at each time. The purple curve overlays the instantaneous emission rate $d\sigma/dt$ for reference; the color bar on the right gives its absolute scale. Top (265 fs): rapid image-charge field build-up on the rising edge suppresses mid-pulse emission. The instantaneous emittance peaks early, reflecting strong

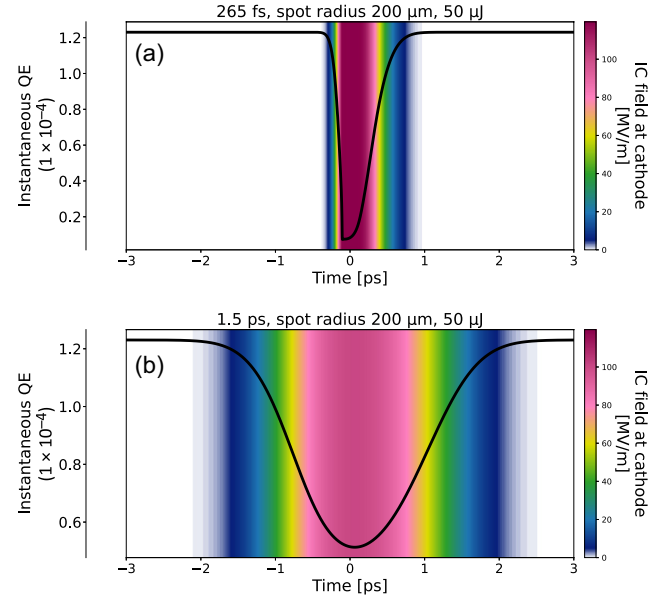


FIG. 7. Temporal evolution of the image-charge (IC) field and instantaneous quantum efficiency for two pulse durations. Color maps: IC field at the cathode $E_{\text{IC}}(0^+, t)$ [MV/m]; black curve: normalized QE(t) from the time-dependent barrier $F_{\text{bar}}(t) = F_{\text{ext}} - E_{\text{IC}}(0^+, t)$. Top (265 fs): rapid IC build-up suppresses QE on the rising edge, followed by slow recovery due to finite charge transport. Bottom (1.5 ps): slower, lower IC build-up yields broader, shallower suppression and weaker asymmetry, illustrating dynamic screening effects vs quasistatic behavior.

Coulomb expansion during the emission burst, followed by a late second emittance peak. Bottom (1.5 ps): emission is spread over a longer interval; earlier accumulation and longer near-surface memory shift the emittance maximum toward the leading side of the pulse, after which acceleration cools the beam and $\epsilon_n(t)$ decays monotonically [38].

This comparison highlights how the collective space-charge dynamics develop differently in short and long pulses, revealing the temporal correlation between charge buildup, field screening, and emittance growth. The results confirm that the onset of space-charge loading and emittance growth is strongly correlated with the temporal profile of the laser pulse, demonstrating how short-pulse emission is self-limited by near-surface collective effects.

We next examine the temporal evolution of the image-charge field and the corresponding instantaneous quantum efficiency for two laser pulse durations, as shown in Fig. 7. The IC field at the cathode, $E_{\text{IC}}(0^+, t)$, exhibits a rapid rise for short pulses, leading to a sharp collapse of the QE on the rising edge as the first emitted electrons quickly build a strong IC field that raises the emission barrier. Recovery on the trailing edge is gradual, limited by the finite transport time of the emitted charge and the falling laser intensity. For longer pulses, the emission occurs over an extended

interval, causing the IC field to build more slowly and reach a lower peak, resulting in a broader, shallower suppression of QE with weaker asymmetry. These dynamics demonstrate that short pulses strongly emphasize barrier modulation and self-screening effects, whereas longer pulses behave nearly quasistatically.

VI. CONCLUSION AND OUTLOOK

The presented model provides a self-consistent description of ultrafast photoemission dynamics, including image-charge buildup and space-charge feedback, validated across a wide range of emission conditions. By resolving the transient evolution of the emission barrier and the instantaneous emission rate, the simulations reproduce key experimental trends such as charge saturation and fluence dependence while also predicting the time-dependent evolution of emittance driven by dynamic field screening. This demonstrates the model's ability to capture both microscopic and collective effects with low computational cost.

Looking forward, this framework will be extended to adaptive spatiotemporal control of the UV laser pulse, enabling real-time adjustment of the emitted charge in the FCC-ee electron source. Such capability is essential for maintaining precise bunch charge control within the continuous top-up injection scheme foreseen [39]. The ultimate goal is to integrate this model into source design, enabling the development of electron sources that meet the requirements of future collider facilities.

DATA AVAILABILITY

The data that support the findings of this article are not publicly available. The data are available from the authors upon reasonable request.

APPENDIX A: MODEL PARAMETERS

This section summarizes the key parameters employed in the photoemission model presented in this work. These parameters define the physical properties, and material-specific constants that govern the emission process and its numerical implementation. Where possible, values are drawn from established literature or experimental measurements to ensure physical consistency and reproducibility. Table I provides an overview of all parameters used, including their symbols, numerical values.

TABLE I. Parameters used in the photoemission model.

Parameter	Symbol	Value	Source
Work function	ϕ	4.5 eV	Literature
Mean transverse energy	MTE	0.25 eV	Assumed
FD prefactor	A_{FD}	3×10^{-5}	Adjustable
Image-charge decay length	z_{MAT}	1.3 μm	Assumed

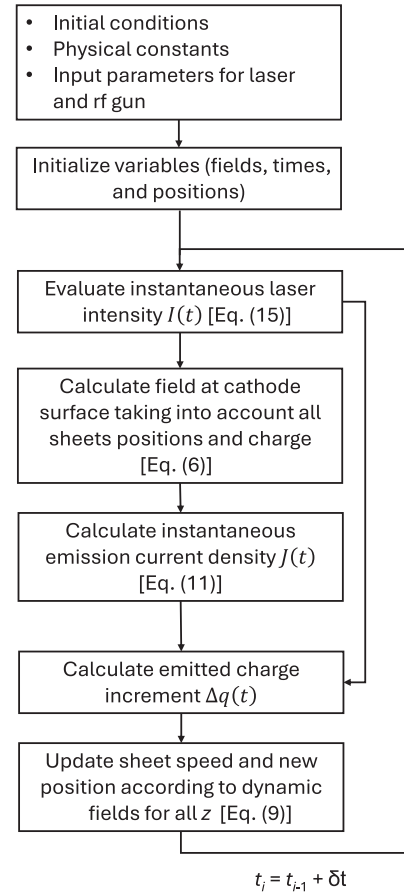


FIG. 8. Flowchart of the time-dependent photoemission model. The algorithm advances the emission process by evaluating the instantaneous laser intensity, computing the effective field at the cathode, determining the emitted charge, and updating the dynamics of emitted sheets.

APPENDIX B: MODEL ALGORITHM

For clarity, the numerical implementation of the model is summarized in the form of a flowchart in Fig. 8. The diagram illustrates the sequence of operations performed at each time step, including field evaluation, emission calculation, and sheet dynamics.

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