Coherent Smith-Purcell Radiation Sources Based on Dielectric Gratings

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We describe a simple method to produce coherent Smith-Purcell radiation with a single dielectric grating. The scheme exploits a strong resonance which enables staggering enhancements of diffraction modes of a subwavelength grating under evanescent wave incidence. The interaction between a continuous electron beam and the resonant mode can bunch electrons, resulting in coherent oscillation, and consequently superradiant Smith-Purcell radiation. Using a higher-order diffraction mode for Smith-Purcell radiation allows one to use low-energy electrons available from tabletop sources. This work potentially paves the way to a compact coherent radiation source, which may find application in the fields of communications, physics, and biology.

In the past decade, growing interest in radiation sources has led to extensive analysis of various interactions between free electrons and surrounding media such as Cherenkov radiation [1–3], transition radiation [4, 5], and Smith-Purcell (SP) radiation [6–13]. Together with the development of densely integrated electron emitters [14, 15], the study of planar periodic structures for SP radiation paves the way toward the realization of on-chip free-electron lasers (FELs) [16]. Dielectric subwavelength gratings (SWGs) have been shown to be especially suitable for efficient spontaneous SP radiation by exploiting strong resonances of the SWG structure [7, 8]. However, so far little work has discussed how to achieve coherent SP radiation with dielectric SWGs, which is desired for applications such as communications, physics, and biology.

Electron bunches interacting with gratings are essential to produce coherent radiation efficiently, where electrons radiate in phase [17, 18]. Electron bunching in a continuous beam is the subject of extensive research recently due to its wide applications including free-electron lasers, ultrafast electron diffraction, and accelerator injectors. Coupling external electromagnetic field to bunch the beam has been widely used thus far in particle accelerators, but suffers from various degrees of complexity [19–22]. Methods to passively bunch the beam are more appealing to a compact SP-FEL. Traditionally, coherent SP radiation sources with metallic gratings adopt two passive bunching approaches which can also apply to dielectric gratings. The Orotron utilizes the interaction between an electron beam and a self-excited standing wave in a Fabry-Perot cavity formed by a grating and a reflection mirror [23–25]. However, the required cavity hinders its extensive application. The superradiant SP-FEL relying on the interaction between an electron beam and a self-excited backward evanescent wave guided by the grating [26–29]. Its capability to generate coherent SP radiation with a single grating has drawn considerable interest, but the power conversion efficiency from electron to SP radiation is rarely optimal since the grating design should also consider the bunching performance. Solutions to this problem usually involve using two cascade gratings [30, 31] or introducing extra cavities [32], but at the cost of complexity.

Here, we present a method in which the strong resonance in a single dielectric SWG is used to achieve passive bunching in tandem with the strong enhancement of the coherent SP radiation. We provide a theoretical analysis of the grating, which leads to physical insights into the resonances excited by the nonradiative incident wave bound at the electron beam. We show that the interaction between an electron beam with the resonant mode can bunch electrons, leading to coherent oscillation, and consequently increasing the field of coherent SP radiation. We also show that the required electron energy can be reduced by using a higher-order diffraction mode for SP radiation. Our method provides an efficient way to produce coherent SP radiation. It may enable a compact on-chip terahertz source and open a door to generate coherent light in all-silicon structures.

We first revisit the physical mechanism of the SP effect. The schematic of a dielectric SWG interacting with a sheet electron beam is shown in Fig. 1(a). The dielectric SWG comprises of alternating high-index ($n_h$) and low-index ($n_l$) bars of thickness $t$. A sheet electron beam moves in close proximity to the SWG at $x = x_0$ with a charge density distribution $q$ per unit length in the $y$ direction and a velocity $v_0 = \beta_0 c$ in the direction along the SWG periodicity. The current density of the electron beam can be written as

$$J(x, z, t) = \frac{q}{2\pi} \delta(x - x_0) \delta(z - v_0 t),$$  \hspace{1cm} (1)

where $k_{z0} = \omega/v_0$. In the free space between the beam and the grating, the field induced by the current can be characterized by its magnetic field:

$$H^{in} = \frac{q}{4\pi} e^{i(k_{z0} z - k_{x0} (x - x_0))},$$  \hspace{1cm} (2)

where $k_{z0} = (k_0^2 - k_{x0}^2)^{1/2}$ is the $x$ wavenumber, $k_0 = \omega/c$ is the free-space wavenumber. With $\beta_0 < 1$, $k_{x0}$ is
FIG. 1. (a) A schematic drawing of a dielectric grating interacting with an electron beam. Electrons (green) travel in the $z$ direction with velocity $v_0$. $n_b$ and $n_a$ are the refractive indices of the high-index bar and low-index bar, respectively. $x = 0$ is located at the grating upper plate. Parameters: $p$, grating period; $t$, grating thickness; $b$, high-index bar width; $a$, low-index bar width; $\theta_n$, radiation angle. (b) $E_z$ field profiles of the 0th diffraction order for electron bunching (up) and the -1st diffraction order for SP radiation (down), with $\theta_{-1} = 90^\circ$. (c) Overview of the electron energy gain as a function of the start phase arising from interaction with the 0th-order diffraction mode.

always imaginary, and thus the field described by Eq. [2] is nonradiative.

Upon evanescent wave incidence, a grating can provide Bloch wavevectors to couple the energy bound at the electron beam into the far field. The diffraction of the evanescent waves accompanying the moving electron yields a series of diffraction modes with $z$ wavenumbers $k_{zn} = k_{z0} + n2\pi/p$, where $n = 0, \pm 1, \pm 2, \text{etc.}$ are the diffraction orders. Rewriting the $z$ wavenumber $k_{zn}$ in terms of the radiation angle $\theta_n$, $k_{zn} = k_0 \cos \theta_n$, we can obtain the well-known relation between the radiation wavelength $\lambda$ and the electron velocity

$$\lambda = \frac{p}{n} \left( \frac{1}{\beta_0} - \cos \theta_n \right). \quad (3)$$

The field of the $n$th diffraction order above the grating with a complex coefficient $r_n$ is

$$H_{n}^{re} = \tilde{y} \frac{r_n q}{4\pi} \text{e}^{ik_0 z_0 \text{e}^{i(k_{zn} z + k_{zn} x)}}, \quad (4)$$

$$E_{n}^{re} = [\tilde{x} k_{zn} - \tilde{z} k_{zn}] \frac{r_n q}{4\pi \omega \varepsilon_0} \text{e}^{ik_{zn} (z + x)} \text{e}^{i(k_{zn} z + k_{zn} x)}, \quad (5)$$

where $k_{xn} = (k_0^2 - k_{zn}^2)^{1/2}$ is the $x$ wavenumber of the $n$th diffraction order. Under evanescent wave incidence, diffraction orders with a real $x$ wavenumber are propagating SP radiation modes, while diffraction orders with an imaginary $x$ wavenumber are evanescent modes that decay exponentially in the $x$ direction. For example, as shown in Fig. [1(b), with $p = \beta_0\lambda$, the -1st diffraction order is propagating while the 0th diffraction order is evanescent. For SP radiation sources with either metallic gratings or dielectric gratings, a large deal of work has been devoted to studying the evanescent-to-propagating wave conversion efficiency $[8, 9, 33, 36]$, while the coexisting evanescent diffraction modes have seldom been utilized.

In contrast, by exploiting a dielectric SWG with a strong resonance, we use an evanescent diffraction mode to bunch electrons, enabling coherent SP radiation. According to Eq. (6), the 0th-order diffraction mode has a $z$ wavenumber $k_{z0}$, indicating that it is phase synchronous with the electrons. Inspired by the velocity bunching scheme in dielectric laser accelerators $[20, 21, 37, 38]$, the longitudinal electric field of the 0th diffraction order can be utilized to accelerate or decelerate electrons. The energy gain of a phase-matched electron over one grating period can be expressed as $\Delta E = -i E_{x0}^{re}(x) e p \sin \omega t_0$, where $E_{x0}^{re}(x) = -k_{z0} r_0 q (4\pi \omega \varepsilon_0)^{-1} \text{e}^{i(k_{x0} z_0 + x)}$, $e$ is the elementary charge, $\omega t_0$ is the start phase of an electron relative to the electric field of the 0th order. As shown in Fig. [2(c)], electrons with $\omega t_0 \in [0, \pi]$ lose energy, and electrons with $\omega t_0 \in [\pi, 2\pi]$ gain energy. Such dependence leads to electron energy modulation, which evolves into density modulation after a distance and thus generates electron bunches. As a result, the incident waves induced by the electron bunches become concentrated in the harmonics of the bunching frequency, thereby leading to coherent oscillation which simultaneously scales up the SP radiation. It should be noted that the reflection coefficient $r_0$ plays an important role in the oscillation. When a constant excitation is stimulated by a continuous electron beam, the 0th diffraction order at different frequencies will be excited. The mode with a high $|r_0|$ builds up energy faster than the rest, eventually dominates and starts oscillation. Besides, to efficiently produce far field radiation, high reflection coefficients are also required for the propagating diffraction modes. We will show that, by making use of a strong resonance, an enhancement of various diffraction modes at the resonant frequency results, including both the propagating modes for SP radiation and the evanescent mode for electron bunching.

To understand dielectric SWGs under evanescent wave incidence, the field profile inside the grating should not be ignored. To this end, we use the waveguide-array (WGA) modes formulation, which was introduced in [39] to explain the reflection of a plane wave impinging on a SWG. Unlike the conventional electromagnetic solvers which compute only the total field, such mode-matching method not only provides a much more straightforward way to extract the diffraction coefficients of an SWG but also leads to intuitive physical insights into its extraordinary properties $[40, 12]$. Below we generalize the analytical method and derive the WGA modes in SWGs under evanescent wave incidence. We will show that the constructive interference of WGA modes at grating interfaces results in a very strong resonance and hence the enhancement of diffraction modes.

Along the $x$ direction, the grating can be treated as
a periodic array of waveguides, where a series of WGA modes exist. Considering the incident wave is TM-polarized, we denote the WGA modes as TM₀, TM₁, TM₂, etc. The lateral magnetic field profile in the z direction for the TMₙ mode can be written as

\[ H_{y_m}(x, z) = A_1 e^{ik_{a,m}x} + A_2 e^{-ik_{a,m}x}, \quad 0 < z < a, \]

\[ H_{y_m}(x, z) = B_1 e^{ik_{b,m}x} + B_2 e^{-ik_{b,m}x}, \quad a < z < b, \]

where the z wavenumbers in the low-index bar and high-index bar are determined by \( k_{a,m} = (n_a^2 k_0^2 - k_{xm}^2)^{1/2} \) and \( k_{b,m} = (n_b^2 k_0^2 - k_{zm}^2)^{1/2} \), respectively, and \( k_{xm} \) is the x wavenumber. Enforcing the continuity of tangential field components at the bar interfaces and the Bloch boundary condition yields the dispersion relation for WGA modes,

\[ 2n_a^2 n_b^2 k_{a,m} k_{b,m} \left[ \cos(k_{a,m}a) \cos(k_{b,m}b) - \cos(k_0p/\beta_0) \right] = (n_a^2 k_{a,m}^2 + n_b^2 k_{b,m}^2) \sin(k_{a,m}a) \sin(k_{b,m}b). \]

Equation (8) suggests that the number of WGA modes can be controlled by designing the dutycycle when the electron energy \( E_0 \) and the ratio of grating period \( p \) to radiation wavelength \( \lambda \) are fixed. For instance, with \( E_0 = 51 \text{ keV}, p/\lambda = \beta_0 = 0.416 (\theta - \phi = 90^\circ), n_b = 3.486, n_a = 1 \), we show the dependence of \( k_x \) on the dutycycle \( b/p \) in Fig. 2(a). As dutycycle increases, more and more modes exist. There is a lower dutycycle limit for each mode, except for the fundamental mode. The dual-mode region where only two modes exist is indicated by the black dash lines.

The introduction of grating boundaries in the x direction confines the WGA modes inside the grating with coupling to the outside diffraction modes. Following the analytical mode-matching method in [39], we can obtain the reflection coefficients \( r_0 \) and \( r_{-1} \). Figures 2(b)–2(d) show that, by choosing appropriate dutycycles and thicknesses such that a resonance is excited, the strong enhancements of the 0th diffraction order as well as the -1st order can be achieved. The highly-ordered patterns in Figs. 2(b) and 2(c) reveal the strong dependence of reflection coefficients on both dutycycle and grating thickness, indicating an interference effect. We will explain this interesting phenomenon by exploring the FP resonance mechanism of those WGA modes.

In an SWG, the field including a finite number of WGA modes propagating downward can be given by \( H_y(x, z) = [H_{y0}(z) \ H_{y1}(z) \ H_{y2}(z) \ldots] e^{ik_xz} [C_0 \ C_1 \ C_2 \ldots]^T \), where \( k_x \) is a diagonal propagation matrix composed of \( k_{xm} \), and \( C = [C_0 \ C_1 \ C_2 \ldots]^T \) is a state vector which characterize the field [39]. After a round trip in the grating, the state vector becomes \( MC \), where \( M = R e^{ik_xL} R e^{ik_xT} \) is the propagation matrix, and \( R \) is the reflection matrix at the interfaces. The original \( M \) obtained by the mode-matching method is undiagonal, indicating that WGA modes are not orthogonal. The diagonalization of the \( M \) matrix provides a new set of orthogonal modes, namely “supermodes” [39]. Accordingly, the phase accumulated by the \( m \)-th order supermode after a round trip corresponds to the phase of the \( m \)-th eigenvalue of \( M \), denoted by \( \phi_m \). Resonances occur when the round-trip phase condition is satisfied by any supermode \( m, \phi_m = 2\pi l, \) with \( l = 0, 1, 2, \) etc. In Fig. 2(e), the combinations of \( t \) and \( b/p \) supporting the resonance dominated by the 1st, 2nd, 3rd-order supermode (denoted SM₁, SM₂, SM₃) are shown by the blue, purple, and green curves, respectively. The patterns in the \( |r_0| \) and \( |r_{-1}| \) maps are grided by different resonant curves, revealing that each supermode contributes differently to those diffraction modes. We notice that a very strong resonance can be obtained when supermodes SM₁ and SM₂ reach their conditions in phase, with \( |\phi_1 - \phi_2| = 2\pi s, \) \( s \) being an even integer. In this case, the strong enhancements of \( |r_0| \) and \( |r_{-1}| \) can be achieved simultaneously, which is the basis of our concept, such as the point marked by a black star in Figs. 2(b) and 2(c).

Now we show that the interaction between an electron beam and the resonant mode of an SWG enables electron bunching and thus coherent SP radiation by using two-dimensional finite-difference-time-domain simu-
FIG. 3. Simulation results of coherent SP radiation sources based on SWG resonances, with electron energy $E_0 = 51$ keV [(a)–(e)] and $E_0 = 4.99$ keV [(f)–(j)]. [(a) and (f)] Snapshots of the particle distribution above the grating. The insets show the development from energy modulation into density modulation of electrons. The numbers of grating period are 30 in (a) and 20 in (f). [(b) and (g)] Phase-space distributions of the electron beam. [(c) and (h)] Field patterns of the $B_y$ component of the total field. [(d) and (i)] Field spectrum of the SP radiation. [(e) and (j)] Power of the upward propagating SP radiation. Parameters: grating period $p = 624$ µm, grating thickness $t = 913$ µm, beam thickness 100 µm.

lations. While our theory is general, we choose a terahertz operation frequency to exemplify our concept. Electron and SWG parameters are in accordance with the black star in Figs. 2(b), with SP radiation along the surface normal direction. Consider an SWG of period $p = 624$ µm, duty cycle $b/p = 0.906$, thickness $t = 913$ µm and number of periods $N = 30$, all placed in an external uniform longitudinal magnetic field of 0.5 T. On top of the grating we place a sheet electron beam of energy $E_0 = 51$ keV, thickness 100 µm, current density $50$ A/cm$^2$. Considering the small transverse extent of the 0th diffraction mode, the distance between the electron beam and the grating is set to be zero. We note that such a single-layer grating could be realized exploiting lithographic and etching techniques[43], and the required electron density can be achieved with state-of-the-art emitters[44].

In Figs. 3(a) and 3(b), we show the particle and phase-space distributions of the electron beam, respectively. Along the longitudinal direction, the electrons interact with the self-excited 0th diffraction order, inducing an energy modulation. After a distance, as shown by the inset in Fig. 3(a), faster electrons intersect slower ones, so the energy modulation develops into density modulation. Owing to the periodicity of the 0th diffraction mode along the $z$ direction ($k_{z0}p = 2\pi$), electron bunches are spaced by one grating period. Due to the evanescent nature of the 0th diffraction mode, electrons with a shorter distance to the grating ($x$ direction) experience stronger energy modulation and thereby form bunches earlier. Figs. 3(c) and 3(d) demonstrate that the SP radiation angle is $90^\circ$ and the radiation frequency is 0.20 THz, in agreement with the values predicted by Eq. (3). Assuming the $y$-direction width of the grating is 1 mm, Fig. 3(e) shows that the power for the SP radiation propagating upwards is 143 W. The power of the downward propagating wave is found to be approximately the same as the power of the upward propagating wave, so the beam-to-wave power conversion efficiency is 11.2%. The preliminary results show that our scheme could not only emit coherent SP radiation, but also outperform a variety of metallic Orotrons[24–26].

The required electron energy is an important metric for free-electron radiation sources. Using electrons with high energies leads to a large size and high cost that severely limit the availability and scalability of the application. Now we show that by utilizing a higher-order diffraction order as the SP radiation mode, our concept is readily to be used for developing low-voltage SP-FELs. For the above SP-FEL example as shown in Fig. 3(a), the synchronicity condition necessitates $p = \beta_0\lambda$, leading to $k_{z0}p = 2\pi$. Equation (8) indicates that, when changing
the normalized electron velocity to $\beta_0/j$ ($j = 2, 3, 4$, etc.),
the solution of the dispersion relation stays the same and
thus the WGA modes and the resonant frequency remain
unchanged. Meanwhile, the $-j$th diffraction order with
$k_z = 0$ becomes a propagating wave along the surface
normal direction. In this case, using the $-j$th diffraction
order for SP radiation and the 0th for bunching allows
us to obtain coherent SP radiation with lower electron
energies. In the following example, we set $j = 3$, with
electron energy $E_0 = 4.99$ keV. The particle and phase-
space distributions of the electron beam are depicted in Figs. 3(f) and 3(g), respectively. With $k_z = 6\pi/p$,
the spacing between electron bunches becomes $p/3$. Fig-
ures 3(h) and 3(i) demonstrate that the field distribution
for the $B_j$ component and the radiation frequency remain
the same as using $E_0 = 51$ keV. Figure 3(j) indicates
that the power for upward propagating SP radiation is
2.8 W, with an efficiency of 2.3%. Compared with using
$E_0 = 51$ keV, the decreasing of power conversion efficiency results from the differences in diffraction coeffi-
cients and the smaller transverse extent of those evanes-
cent modes.

In conclusion, we presented the concept of using a sin-
gle SWG for coherent free-electron radiation sources. We
showed that a strong resonance could be obtained in an
SWG by engineering the interference of WGA modes.

The interaction between a continuous electron beam and
the resonant mode of an SWG can lead to oscillation,

enabling coherent SP radiation. By using a higher-order
diffraction mode for SP radiation, the required electron
energy can be reduced. As an outlook, we note that our
study suggests that similar devices could be realized with
metallic structures by making use of the special SP radia-
tion from an open resonator array [45]. In addition, based
on the scaling law of Maxwell’s equations, our design is
readily transferrable to optical frequencies, potentially
achieving an efficient silicon light source, which might
find broad applicability in industries such as communi-
cations, sensing and lighting. With rapid advances in the
development of FEs, our study enables a viable way of
realizing a low-voltage on-chip coherent radiation source
that might change the affordability of FELs.

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