A directional, ultrafast and integrated few-photon source utilizing the interaction of electron beams and plasmonic nanoantennas

Nahid Talebi
Max Planck Institute for Intelligent Systems, Heisenbergstr. 3, 70569 Stuttgart, Germany
E-mail: talebi@is.mpg.de

Received 15 January 2014, revised 18 March 2014
Accepted for publication 9 April 2014
Published 9 May 2014
New Journal of Physics 16 (2014) 053021
doi:10.1088/1367-2630/16/5/053021

Abstract
The possibility of using plasmonic nanoantennas in interaction with electron beams to form a wideband few-photon source is demonstrated here using numerical investigations. The generated photons are guided inside a waveguide, which is coupled to the nanoantennas. Using two replicas of a nanoantenna at an optimized distance, and providing a guided path for the emitted photon from the first nanoantenna to the second one by positioning a resonator feedback element, it is possible to form an interference between the photons generated by the two nanoantennas. The interference phenomenon is used to produce a unidirectional flow of power in the waveguide, as well as a high collection efficiency. The investigations are carried out in a self-consistent way, utilizing a conjugate Maxwell and Lorentz system of equations, in order to be able to simulate the modulation of the electron velocity due to the interaction with an optical system which is initially at rest.

Online supplementary data available from stacks.iop.org/NJP/16/053021/mmedia

Keywords: few-photon source, electron energy loss probability, photon generation probability, cathodoluminescence, poynting theorem, interference, plasmon
1. Introduction

Ultrafast science, intended to investigate the time-resolved response of matter and structures, underwent tremendous success in the previous decade due to the introduction of ultrafast electromagnetic sources: so-called free-electron lasers (FELs) [1–3]. In such a configuration, a relativistic electron beam travelling along the axis of the magnetic gratings is utilized as a laser-like radiation source caused by the self-amplified spontaneous emission. The FEL emits in the extreme ultraviolet (EUV) range of the electromagnetic spectrum, which was a challenge to reach prior to that [1]. Besides the EUV FELs, the possibility of tuning the emitted photon energies from millimeter waves [4, 5] to x-rays [6] increases the applicability of FEL in a variety of research areas. In fact the parameters being considered in the design of FELs are the brightness, the generated photon frequency, the temporal pulse duration and the coherence. In addition to the mentioned parameters, there has been also a tremendous attempt to decrease the size of photon sources to the micrometer range [7, 8].

Besides the Larmor mechanism of radiation utilized in the design of FELs, other mechanisms of radiation, such as the Smith–Purcell effect [9–11] and transition radiation, can also be used to introduce novel optical sources. For example, a nanoscale optical source, composed of a well milled into a stack of gold and silica layers, has been demonstrated, showing approximately 20 nm broadening at the emission output [12, 13]. The interaction of electron beams with a planar metamaterial has been elsewhere demonstrated as a coherent source of optical radiation [14].

In addition to the trend in ultrafast optics, the control over the emission process of the optical source aiming at single-photon generation has prospective applications in quantum computing [15] and fundamental quantum optics experiments [16]. Spontaneous emission initially offers a choice for achieving the single-photon limit, though incorporation of other phenomena like electroluminescence is also possible [17]. Spontaneous emission from quantum dots [18–20], organic molecules [21, 22] and defect centers like nitrogen vacancies [23, 24] has so far been considered as possible single-photon sources. In order to achieve more directionality and hence a higher collection efficiency, the coupling of the emitters to higher-index media [25–27], nano-antennas [19, 28–30], nanowires [18] and photonic crystals [31] has been proposed. Besides the previous processes for controlling the directionality of the emitted single photons, the interaction of a single electron beam with a thin optical fiber has recently been proposed as an efficient method for generating a nearly deterministic single photon in an optical fiber [32]. Moreover, the possibility of coupling the electron beams to the nitrogen-vacancy center as a single photon emitter has been also experimentally demonstrated, as a way to control the generation of single-photon states [33].

The single-photon sources investigated so far do not guarantee the propagation of the generated photons inside an integrated waveguide, but the investigated photon sources sustain a directional radiation pattern into the surrounding medium. However, for applications in quantum computing, it is highly demanded that the emitted photons can be controlled and guided in optical integrated circuits on the nanometer-scale. Here, it will be shown that a hybrid circuit of nanoantennas coupled to plasmonic waveguides can be used for steering the free-electron beams and controlling their emission to single- and few-photon regimes. The advantage of introducing the integrated on-chip structure as a few-photon source, beside following the trend of miniaturization, is the ability to couple the waveguide to other processing units, such as resonators and filters. More importantly, the proposed circuit can be utilized to
steer the emission process of the electron in the desired direction, without the need to introduce external electromagnetic excitations. Such external excitations have the problem of imperfect mutual coherence of the particle and optical excitations, which violates the efficiency of the method due to the time-jitter problem for the electron beams [34, 35].

In order to investigate the functionality of the proposed circuit and design the necessary parameters accordingly, a home-built particle-in-cell numerical code within the context of a finite differentiation method is used. The accuracy and efficiency of the introduced finite-difference time-domain code with an embedded electron source has been shown elsewhere within the limits of an undepleted pump approximation, which neglects the recoil experienced by the electrons during the interaction [36]. Here, in order to compute the recoil experienced by the electron beam in the time of the interaction, Maxwell’s equations are coupled to the relativistic Lorentz equation. The resulting system of equations is solved using a self-consistent finite differentiation scheme.

The photon generation (PG) probability is analytically computed using the integration of the Poynting vector into the far-field in the frequency domain [37], which is valid within the time-harmonic field theorem [38, 39]. Since the generated plasmons are drastically wideband in the frequency domain, the question arises how the emitted power from the structure can be analytically described by a time-harmonic Poynting vector integral. In order to approach this question, one should address the definition of the power spectrum detected by the detectors. While the spontaneous power is clearly defined by the time-domain Poynting theorem, the detected averaged power spectrum due to the electromagnetic wide-band pulses will be shown to be efficiently describable by assuming a periodic excitation. This is a justified assumption, since the laser sources in optical experiments or electron sources in transmission electron microscopes sustain a repetition rate of the order of GHz.

2. Proposed structure and first principles

A swift electron moving in free space cannot emit photons due to the energy–momentum mismatch between the two particles. However, a particle can provide enough momentum in the near-field for an inelastic interaction with the electron, so that an electron traveling close enough to the structure can excite the sample to a higher photonic state. The structure will then emit a photon and relax to the vacuum state.

The elementary processes of the inelastic interaction of electrons and any optical system will lead to a loss of energy from the electron. Electron energy-loss spectroscopy (EELS) can measure the probability of the electron to lose a certain amount of energy, due to the excitation of single or multiple quanta of photonic states. However, the PG spectrum or the so-called cathodoluminescence shows the probability of the system to re-emit the absorbed quanta of photons. Relations between these two elementary and secondary processes can provide us with information about the relaxation and loss mechanisms occurring in the system.

Classically, EELS and PG spectra are given by [37]:

$$
\Gamma_{\text{PG}}^{\text{PG}}(\omega) = \frac{1}{\hbar \omega} \iint_{s} \text{Re} \left\{ \vec{E}(\vec{r}, \omega) \times \vec{H}^{*}(\vec{r}, \omega) \right\}. ds
$$

(1)
\[ \Gamma^{EELS}(\omega) = \frac{-1}{\hbar \omega} \iiint_v \text{Re} \left\{ \tilde{E}(\vec{r}, \omega) \cdot \tilde{J}^{*}(\vec{r}, \omega) \right\} \, dv \]  

(2)

in which \( \hbar \) is the Planck constant, \( \tilde{E}(\vec{r}, \omega) \), \( \tilde{H}(\vec{r}, \omega) \), and \( \tilde{J}(\vec{r}, \omega) \) are the electric field vector, the magnetic field vector and the current density function, respectively. Deriving an energy conservation criterion using the Poynting theorem as shown in the appendix, it can be shown that these two elementary and secondary mechanisms are related by the energy dissipation spectra caused by the anti-Hermitian parts of the permittivity \( \tilde{\varepsilon}(\omega) \) and permeability \( \tilde{\mu}(\omega) \) tensors, as:

\[ \Gamma^{EELS}(\omega) = \Gamma^{PG}(\omega) \]

\[ - \frac{1}{\hbar} \iiint_v \tilde{H}(\vec{r}, \omega) \cdot \text{Im} \left\{ \tilde{\varepsilon}_{AH}(\omega) \right\} \cdot \tilde{H}^{*}(\vec{r}, \omega) \, dv \]

\[ - \frac{1}{\hbar} \iiint_v \tilde{E}(\vec{r}, \omega) \cdot \text{Im} \left\{ \tilde{\mu}_{AH}(\omega) \right\} \cdot \tilde{E}^{*}(\vec{r}, \omega) \, dv \]  

(3)

In this paper, we consider a plasmonic circuit for the interaction medium. The advantages of using such a configuration for engineering the local density of states over the dielectric circuits are due to the size and the huge near-field enhancement, which equivalently enhances the probability of spontaneous emission in interactions with the emitters. However, to incorporate the emitted photons in quantum computing, it would be best to couple the radiation pattern into a plasmonic waveguide, which on the other hand can be used as a sub-diffraction medium for guiding the emitted photon.

For the configuration we propose here (figures 1(a) and (b)), we use single or double nanoantennas coupled to a rib plasmon waveguide, positioned on a silica substrate. The thickness of the silica layer is taken to be 300 nm throughout the paper, and it is positioned on a silicon layer of infinite thickness. Structure (a) is composed of a nanoantenna with a thickness of \( H_M = 30 \) nm, a width of \( W = 30 \) nm and a length of \( L_y = 147 \) nm, coupled to the rib waveguide mentioned above. The mentioned nanoantenna has a first-order resonance at the energy of \( E = 1.34 \) eV, without considering its coupling effect to the rib waveguide. The gap between the nanoantennas and the rib waveguide is taken as \( d_1 = 20 \) nm. The electron impact parameter and energy are taken as \( b = 2 \) nm and \( U_e = 200 \) keV, respectively. Numerical investigations show that a thinner rib waveguide with \( W \leq 50 \) nm is preferable, since a thick waveguide sustains higher material losses for the photons propagating at energies above 1.5 eV. Moreover, in order to provide an efficient coupling of the nanoantennas to the fundamental mode of the waveguide, a waveguide with a width less than \( W = 40 \) nm should be used.

There exist two mechanisms of radiation from the free-electron beam during their interaction with the above-mentioned structures: namely Cherenkov radiation, due to the interaction of the incident electrons with the substrate (silica), and also transition radiation. These mechanisms result in the excitation of several photons in the coupled system of the antenna and the rib waveguide. Since the aim is to design a few-photon source, Cherenkov radiation should be excluded in the desired frequency range. This is the reason for the choice of a silica substrate here. The EELS spectra due to Cherenkov radiation can be computed analytically [40]. Figure 1(c) shows the computed EELS spectra for an electron traveling parallel to the interfaces of silica/air and gold/air, at a distance of 1 nm. It is evident that the
contribution of Cherenkov radiation from the silica substrate at energies below 9 eV is negligible. This means that transition radiation from the plasmonic system is the only mechanism for the generation of photons below 10 eV.

In order to simulate the dynamics of the electron and plasmon interactions, a numerical code based on the combined Maxwell and Lorentz equations has been developed, within the finite differentiation algorithm. The efficiency of our previously developed code for computing EELS spectra within the limits of the undeleted pump approximation has already been presented [36]. Here, we include the Lorentz equation to simulate the velocity modulation and dynamics of electron beams and to be able to calculate the Larmor radiation of the electrons in a self-consistent approach. In order to model the dispersive permittivity of the introduced materials, a Drude model with two added critical-point functions is used [41, 42]. Previous numerical results using this model show perfect agreement with the experimental results, which demonstrates the suitability of this model for time-domain analysis [43, 44].
3. Results and discussion

First, the structure shown in figure 1(a) is considered. Figure 2(a) shows the computed EELS spectrum for a 200 keV electron moving adjacent to the nanoantenna, with an impact parameter \( b = 2 \) nm. Figure 2(b) shows the guided and scattered averaged power spectra for the same structure, \( P_{\text{scat}} = P_y^- + P_z^+ + P_z^- \), where \( P_y^-, P_z^+, P_z^- \) are the power flows in the \(-y, -z, +z\) directions. (c) Relative electron velocity \( \beta (t) = V_e (t) / c \) versus time, along the \( x-, y-, \) and \( z\)-directions. The shadowed region shows the velocity of the electron while passing the nanorod antenna. (d) The induced \( x\)-component of the electric field along the trajectory of the electron, versus time. White and red dashed lines show the light lines in air and silica, respectively. Cherenkov radiation of the electron along the trajectory is apparent in the region between the two optical lines, while a clear plasmon resonance launched by the nanoantenna oscillates in time with a dephasing of approximately 25 fs.

Figure 2. (a) The computed EELS spectrum for a single gold nanorod interacting with a rib waveguide. The inset shows the induced \( z\)-component of the electric field at an energy of 1.3 eV. (b) The computed guided and scattered averaged power spectra for the same structure. \( P_{\text{scat}} = P_y^- + P_z^+ + P_z^- \), where \( P_y^-, P_z^+, P_z^- \) are the power flows in the \(-y, -z, +z\) directions. (c) Relative electron velocity \( \beta (t) = V_e (t) / c \) versus time, along the \( x-, y-, \) and \( z\)-directions. The shadowed region shows the velocity of the electron while passing the nanorod antenna. (d) The induced \( x\)-component of the electric field along the trajectory of the electron, versus time. White and red dashed lines show the light lines in air and silica, respectively. Cherenkov radiation of the electron along the trajectory is apparent in the region between the two optical lines, while a clear plasmon resonance launched by the nanoantenna oscillates in time with a dephasing of approximately 25 fs.
considered, to ensure that only far-field radiation contributes to the calculated power. More investigations show that $P_y$ has the highest contribution to the scattered power. This can be understood by considering the recoil that the electron receives along the different directions. Figure 2(c) shows the velocity of the electron versus time along the x-, y- and z-directions. The electron velocity gradually decreases with a constant rate due to Cherenkov radiation, until it reaches the nanoantenna. By approaching the nanoantenna, the electron receives recoil due to the induced charges at the nanoantenna, resulting in a marked change of the momentum of the electron along the x- and y-directions, while the deceleration along the z-direction remains almost constant. It is apparent that the change in velocity of the electron along z is mostly caused by Cherenkov radiation, rather than by transition radiation.

Figure 2(d) shows the magnitude of the x-component of the induced electric field along the trajectory of the electron versus time. Cherenkov radiation causes a bright light line with an energy–momentum relation of $\omega = \omega_0 / V$, where $V$ is the electron velocity. The plasmon resonance launched by the nanowire oscillates with a dephasing time of approximately 25 fs.

In order to control the dephasing of the plasmon resonances and the directionality of the emission, the structure depicted in figure 1(b) can be used, which includes two nanoantenna at a separation $P$.

The time for the electron to travel between two nanoantennas is $\Delta t = P / V$ in the laboratory frame. If the travel time for the electron is set equal to half of the time-period of the plasmonic resonance, i.e. $\Delta t_{ph} = \pi \hbar / E_{ph}$, it is possible to match the emitted photons from both nanoantennas in a constructive way inside the waveguide. In order to satisfy this criterion for the photon energy of $E_{ph} = 1.32 \text{ eV}$, the distance between the nanoantennas should be $\Delta x = 327 \text{ nm}$. The computed average power spectra for the guided and radiated waves are shown in figure 3(a). The photon generation probability is almost doubled in comparison with the previous structures, which demonstrates the emission of two photons due to the transition radiation from two nanoantennas located along the trajectory of the electron. The comparison of the power spectra of figure 3(a) with figure 2(b) illustrates the higher probability of the generated photons to propagate inside the waveguide, which is caused by the constructive interference of the emitted waves from the nanoantennas.

Figure 3(b) shows the magnitude of the x-component of the electric field along the trajectory of the electron. Comparing this figure with figure 2(d), it is obvious that two plasmon resonances with perfect time matching oscillate in time, each lasting for approximately 25 fs. Indeed, with such a configuration, it is possible to demonstrate the constructive interference of the photons emitted from the interaction of an electron and a plasmonic nanoantenna with the photons generated with another replica of the structure, regardless of the time-jitter caused by the electron optics or electron sources, considering that only one electron interacts with the structure at a given time. It is interesting to note that this interference is quite visible in the scattered power spectra of figure 3(a), compared with figure 2(b).

Figure 3(c) shows the field profile for the z-component of the electric field in the x–y plane located 4 nm above the upper surface of the gold nanostructure, at different times of $t_1 = 10 \text{ fs}$, $t_2 = t_1 + T / 8$ and $t_3 = t_1 + T / 4$. At $t_1$, the photons are localized in the region between the nanoantennas, and they are released after a quarter of the time-period of the plasmon resonance. The emission of the photons from the rib waveguide is in a bi-directional pattern, propagating in both the $+x$ and $-x$ directions.
The possibility of using a configuration of coupled nanoantennas in interaction with the wires to control the directionality of the induced plasmons in the wire has been studied elsewhere [45] for the case of optical plane-wave illuminations. This technique is called phase-engineering, due to the algorithm that is utilized to compensate for the phase lag between the incident mechanism and the fundamental mode of the rib waveguide. The same technique can be used here for coupling the photons emitted from the nanoantennas to the waveguide, and to control their propagation in the waveguide in a single direction. Considering the energy of 1.3 eV for the generated photons, the necessary distance between the nanoantennas is set as $\Delta x = 130$ nm, due to the effective wavelength of the guided plasmons inside the rib-waveguide at this energy. The averaged power spectra for the guided and scattered powers are shown in figure 4(a). Although a more pronounced peak for the emission probability can be achieved at the energy of 1.32 eV, the overall emission probability and collection efficiency of the photons

Figure 3. (a) The computed averaged guided and scattered power spectra for an electron interacting with the structure depicted in figure 3(a), with $\Delta x = 327$ nm. (b) The magnitude of the induced $x$-component of the electric field along the trajectory of the electron, versus time. The red and white dashed lines show the light line inside the silica and air. It is visible that the plasmonic oscillations caused by both nanoantennas are constructive in time, giving rise to a long-lasting resonance in both nanoantennas. The emission caused by these two nanoantennas propagates in both directions in the rib waveguide. (c) The field profile for the induced $z$-component of the electric field at an energy of $E_{ph} = 1.3$ eV, in the $x$–$y$ plane 4 nm above the upper surface of the structure, at different times: $t_1 = 10$ fs, $t_2 = t_1 + T/8$ and $t_3 = t_1 + T/4$, where $T = 2\pi\hbar/E_{ph}$. 

The possibility of using a configuration of coupled nanoantennas in interaction with the wires to control the directionality of the induced plasmons in the wire has been studied elsewhere [45] for the case of optical plane-wave illuminations. This technique is called phase-engineering, due to the algorithm that is utilized to compensate for the phase lag between the incident mechanism and the fundamental mode of the rib waveguide. The same technique can be used here for coupling the photons emitted from the nanoantennas to the waveguide, and to control their propagation in the waveguide in a single direction. Considering the energy of 1.3 eV for the generated photons, the necessary distance between the nanoantennas is set as $\Delta x = 130$ nm, due to the effective wavelength of the guided plasmons inside the rib-waveguide at this energy. The averaged power spectra for the guided and scattered powers are shown in figure 4(a). Although a more pronounced peak for the emission probability can be achieved at the energy of 1.32 eV, the overall emission probability and collection efficiency of the photons
is lower than in the case of two nanoantennas with the distance of $\Delta x = 327$ nm. This can be understood by investigating the time structure of the plasmonic resonances, launched by both nanoantennas. Figure 4(b) shows the $x$-component of the electric field along the trajectory of the electron versus time. The red and white dashed lines show the light lines inside silica and air. The plasmonic oscillations caused by both nanoantennas are constructive in time, giving rise to a short-lasting resonance in both nanoantennas, along with a unidirectional emission inside the rib-waveguide. Figure 4(c) shows the field profile for the induced $z$-component of the electric field at an energy of $E_{ph} = 1.3$ eV, in the $x$–$y$ plane 4 nm above the upper surface of the structure at different times: $t_1 = 10$ fs, $t_2 = t_1 + T/8$ and $t_3 = t_1 + T/4$, where $T = 2\pi\hbar/E_{ph}$. The induced photons in the waveguide propagate along the $+x$-direction.
The directionality of the launched plasmon in the wire, which is computed as the ratio of the power of the wave propagating along the $+x$-direction to the one propagating along the $-x$-direction, is approximately 4.7 for the mentioned structure at the photon energy of 1.3 eV.

The computed velocities of the electron versus time for the structures with $\Delta x = 130$ nm and $\Delta x = 327$ nm are shown in figure 5(a). It is visible from this figure that the electron receives two recoils along its trajectory, with similar time structures. When the electron approaches a nanoantenna it is initially repelled, after which it becomes attracted when it moves away from the nanoantenna at a scattering angle of $\theta = 1.5725 \times 10^{-8}$ rad, as depicted in figure 5(b). For two nanoantennas positioned at a distance of $\Delta x = 327$ nm, which is not within the near-field region, the electron is scattered again with the same scattering angle from the second nanoantenna, resulting in an overall scattering angle of $2\theta$ from its original trajectory. However,
the overall scattering angle for the case of the nanoantennas positioned at a distance of \( \Delta x = 130 \text{ nm} \) is nearly the same as the scattering angle from a single nanoantenna.

Because the power guided and reflected by the structures is approximately equal, the question is how one can optimize the structure to gain more directionality and collection efficiency. Since a large amount of power is reflected into the \( y \)-component of the Poynting vector, the introduction of a feedback element to redirect the power flowing into this component into the wire can greatly enhance the collection efficiency of the emitted photons. For the case of the feedback element, a split-ring resonator is chosen, as depicted in the inset of figure 6(a), with \( L_v = 110 \text{ nm} \) and \( L_h = 100 \text{ nm} \). The first reason for this choice is that the local field distribution along the trajectory is the same for both the nanoantenna and the vertical arms of the split-ring resonator, which provides a perfect coupling for the photons produced in both elements. The second reason is to provide a guided path for the generated photons from the first interaction point to the second interaction point, which, on the other hand, can provide a means for constructive interference. This interference phenomenon can be seen from the computed EELS spectra, depicted in figure 6(a). For the geometry of the structure to be designed perfectly, the system should sustain the same resonance at 1.3 eV, while the Fano-shaped resonance which was obvious previously for the structure without feedback, is changed to another resonance peak, due to the strong interference of the generated photons at the two interaction points.

In order to investigate the collection efficiency of the circuit with the feedback element in comparison with the previous structure, the ratio of the guided power to the scattered power is computed. It should be noted that since both power spectra are computed at the same frequencies, the mentioned ratio is equal to the probability for the generated photons to be guided inside the waveguide, over the probability for the photons to be scattered away. This ratio is shown in figure 6(b) for the circuit with a feedback element and without a feedback element.

The power spectra for the guided and scattered photons are also depicted in the inset. This figure should be compared with figure 4(a), which demonstrates a strong reduction of the scattered power at the desired photon energy of 1.32 eV. Figure 6(c) shows the computed velocity of the electron versus the time-of-flight for the electron, through the simulation domain. While there is no significant change for the \( x \)- and \( z \)-components of the velocity, the \( y \)-component of the velocity is changed into a rather monotonic reduction. The computed \( x \)-component of the induced electric field along the trajectory of the electron is depicted in figure 6(d). A strong modulation of the generated optical field due to the presence of the resonator is visible, which results in a generation of a series of optical pulses with a high damping rate. This modulation shows the efficiency of the feedback element, which results in the re-excitation of the damped plasmon resonances inside the nanoantennas, providing a channel between the two interaction points.

Figure 6(e) shows the \( z \)-component of the electric field versus time, within half of the time period of the excited plasmon. It is visible that the excited plasmon sustains a unidirectional flow of power along the +\textit{x}-direction. The probability for the photon generated at the energy of 1.3 eV to propagate to the right side of the wire is 4.6 times the probability for the same photon to propagate to the left. It is important to note that there is no significant difference for this ratio in comparison with the structure without the feedback element. In fact, the introduction of the feedback element has resulted in an overall increase in the collection efficiency of the generated
Figure 6. (a) The EELS spectra for an electron moving adjacent to the structure with and without a split-ring resonator as a feedback element. The inset shows the structure, including a split-ring resonator, the electron trajectory through the gap between the nanoantennas and the feedback element. (b) The ratio of the photons propagating inside the waveguide and the scattered photons for the structures with and without a feedback element. The guided and scattered average power spectra are depicted in the inset. (c) The velocity of the electron versus time projected along the $x$- and $y$-axes. (d) The magnitude of the $x$- and $y$-components of the induced electric field along the trajectory of the electron. (e) The field profile for the induced $z$-component of the electric field at an energy of $E_{ph} = 1.3$ eV in the $x$–$y$ plane 4 nm above the upper surface of the structure, at different times: $t_1 = 5$ fs, $t_2 = t_1 + 0.1$ $T$, $t_3 = t_1 + 0.2$ $T$, $t_4 = t_1 + 0.3$ $T$, and $t_5 = t_1 + 0.4$ $T$, where $T = 2\pi\hbar/E_{ph}$ (see the movies in the supplementary data).
photons, by decreasing the probability of the photons that are scattered into the surrounding medium.

In fact, the unidirectional flow of power in the waveguide is achieved by incorporating two interaction points for the electrons and plasmons, along with providing a path for the generated photons to interfere in a constructive or destructive way. It is instructive to compare the mentioned situation with the case of an electron traveling normal to the plane of the circuit, in which only a single interaction point exists. Comparing these two cases, as demonstrated in the supplementary data, is an illustrative way to study the differences between the coupling and interference phenomena. Supplementary figure 1 compares these two cases according to the induced electric field projected along the trajectory of the electron, which shows a strong modulation of the field due to the interference of the plasmons launched at two interaction points. This interference results in the extinguishment of the power flow in one direction inside the waveguide. The movies provided as a supplement to this paper show the $z$-component of the induced electric field 3 fs after the interaction at a time period of the excited photons, for both cases of the electron trajectories parallel and normal to the SiO$_2$/air interface (available from stacks.iop.org/NJP/16/053021/mmedia).

Moreover, as an outlook for further experimental investigations, the coupling of the system to single photon emitters, such as nitrogen vacancy centers, can be considered. In such a configuration, the vacancy centers are positioned near to the nanoantenna emitters. The relative excitation time of single photons at the first and second interaction points can be precisely controlled by the velocity of the electron, in order to achieve the phase-matching condition for a unidirectional flow inside the waveguide.

4. Conclusion

A combined system of gold nanoantennas, a waveguide and a split-ring resonator has been investigated here for its possibility to be used as a circuit for single- and two-photon sources in interaction with a relativistic swift electron beam. This system sustains a high collection efficiency and can be designed for a unidirectional flow of power in the system. More importantly, it is in the form of an integrated circuit, which can be realized experimentally using well-known lithography techniques. The unidirectional flow of power in the introduced waveguide, which has been achieved by means of a phase-matching technique, is of critical importance for any processing unit to be followed after the waveguide, for example for quantum computing applications.

Acknowledgement

The author gratefully acknowledges the Alexander von Humboldt Foundation for financial support, and Prof. Peter van Aken, Dr. Ralf Vogelgesang and Dr. Wilfried Sigle for helpful discussions.

Appendix: EELS and PG spectra

In order to derive a law for the conservation of energy in the investigated electromagnetic system, Poynting’s theorem should be applied. We start with the Maxwell equations given by
\[ \nabla \times \vec{E}(\vec{r}, t) = -\frac{\partial \vec{B}(\vec{r}, t)}{\partial t} \] and \[ \nabla \times \vec{H}(\vec{r}, t) = \frac{\partial \vec{D}(\vec{r}, t)}{\partial t} + \vec{J}(\vec{r}, t) \], in which \( \vec{E}(\vec{r}, t) \), \( \vec{H}(\vec{r}, t) \), \( \vec{D}(\vec{r}, t) \) and \( \vec{B}(\vec{r}, t) \) are the electric field, magnetic field, electric displacement and magnetic flux density, respectively, computed at the displacement vector \( \vec{r} \) and time \( t \). \( \vec{J}(\vec{r}, t) \) is the current density function. We will consider the case of a dispersive medium with the constitutive relations as below:

\[ \vec{D}(\vec{r}, t) = \int_{-\infty}^{+\infty} \hat{\varepsilon}(\vec{r}, t - \tau) \vec{E}(\vec{r}, \tau) d\tau \]
\[ \vec{B}(\vec{r}, t) = \int_{-\infty}^{+\infty} \hat{\mu}(\vec{r}, t - \tau) \vec{H}(\vec{r}, \tau) d\tau \]  
(A1)

where \( \hat{\varepsilon}(\vec{r}, t) \) and \( \hat{\mu}(\vec{r}, t) \) are the time domain electric permittivity and magnetic permeability tensors, which can be obtained from the local dispersion relations of the materials using the Fourier transform as below:

\[ \hat{\varepsilon}(\vec{r}, t) = \frac{1}{2\pi} \int_{-\infty}^{+\infty} \hat{\varepsilon}(\vec{r}, \omega) e^{i\omega t} d\omega \]
\[ \hat{\mu}(\vec{r}, t) = \frac{1}{2\pi} \int_{-\infty}^{+\infty} \hat{\mu}(\vec{r}, \omega) e^{i\omega t} d\omega \]  
(A2)

Using the usual procedure for deriving Poynting’s theorem from Maxwell’s equations, the following relation will be obtained:

\[ \nabla \cdot \left( \vec{E}(\vec{r}, t) \times \vec{H}(\vec{r}, t) \right) + \vec{H}(\vec{r}, t) \cdot \frac{\partial \vec{B}(\vec{r}, t)}{\partial t} + \vec{E}(\vec{r}, t) \cdot \frac{\partial \vec{D}(\vec{r}, t)}{\partial t} = -\vec{E}(\vec{r}, t) \cdot \vec{J}(\vec{r}, t) \]  
(A3)

In order to attribute the spontaneous electromagnetic power to the power spectrum measured using the spectroscopy techniques, we assume the field components as periodic functions of time. It is a justified assumption, since the output of all the laser sources utilized in the experiments are periodic over a specified rate, e.g. considering the terahertz modulation of the femtosecond emitted pulses. Since we are dealing with periodic field components, we can expand all the field components with a Fourier series as below:

\[ \vec{A}(\vec{r}, t) = \sum_{n=-\infty}^{+\infty} \vec{A}_n(\vec{r}) \exp\left(\frac{2\pi n}{T} t\right) \]  
(A4)

in which \( \vec{A}(\vec{r}, t) \) are the vector functions describing the field components and \( T \) is the temporal periodicity of the incident pulses. Note that the spatial vector components can be obtained using \( \vec{A}_n(\vec{r}) = \frac{1}{T} \int_{-T/2}^{+T/2} \vec{A}(\vec{r}, t) \exp\left(-i\omega_n t\right) dt \), where we introduced the sampling frequencies as \( \omega_n = 2\pi n/T \). Using the introduced notations and using the real valued-ness of the electromagnetic fields as \( \vec{E}_n(\vec{r}) = \vec{E}_{-n}^*(\vec{r}) \), equation (A3) can be written as:
\[
\mathbf{\nabla} \cdot \left( \sum_{m=-\infty}^{+\infty} \sum_{n=-\infty}^{+\infty} \tilde{E}_n^* (\mathbf{r}) \times \tilde{H}_m (\mathbf{r}) \exp \left( i (\omega_m - \omega_n) t \right) \right) + \sum_{m=-\infty}^{+\infty} \sum_{n=-\infty}^{+\infty} \tilde{H}_n^* (\mathbf{r}) \cdot i \omega_m \tilde{\mu} (\omega_m) \cdot \tilde{H}_m (\mathbf{r}) \exp \left( i (\omega_m - \omega_n) t \right) + \sum_{m=-\infty}^{+\infty} \sum_{n=-\infty}^{+\infty} \tilde{E}_n^* (\mathbf{r}) \cdot i \omega_m \tilde{\varepsilon} (\omega_m) \cdot \tilde{E}_m (\mathbf{r}) \exp \left( i (\omega_m - \omega_n) t \right)
\]

\[
= - \sum_{m=-\infty}^{+\infty} \sum_{n=-\infty}^{+\infty} \tilde{E}_n^* (\mathbf{r}) \cdot \tilde{J}_m (\mathbf{r}) \exp \left( i (\omega_m - \omega_n) t \right) \tag{A5}
\]

The second and third terms are related to the stored magnetic and electric energy densities, respectively. In a similar approach to [38], one can divide the second term in equation (A5) into two parts, in one changing the indices \( m \rightarrow -n \) and \( n \rightarrow -m \) to obtain:

\[
U_m = \frac{1}{2} \sum_{n=-\infty}^{+\infty} \sum_{m=-\infty}^{+\infty} \tilde{H}_n^* (\mathbf{r}) \cdot i \omega_m \tilde{\mu} (\omega_m) \cdot \tilde{H}_m (\mathbf{r}) \exp \left( i (\omega_m - \omega_n) t \right)
\]

\[
+ \frac{1}{2} \sum_{n=-\infty}^{+\infty} \sum_{m=-\infty}^{+\infty} \tilde{H}_n^* (\mathbf{r}) \cdot \left( - i \omega_m \tilde{\mu}^* (\omega_m) \right) \cdot \tilde{H}_m^* (\mathbf{r}) \exp \left( i (\omega_m - \omega_n) t \right)
\]

\[
= \frac{1}{2} \sum_{n=-\infty}^{+\infty} \sum_{m=-\infty}^{+\infty} \tilde{H}_n^* (\mathbf{r}) \cdot \left( i \omega_m \tilde{\mu} (\omega_m) - i \omega_m \tilde{\mu}^* (\omega_m) \right) \cdot \tilde{H}_m (\mathbf{r}) \exp \left( i (\omega_m - \omega_n) t \right) \tag{A6}
\]

The term in the bracket on the right-hand side of equation (A6) implies the decomposition of the permeability tensor into two parts, Hermitian and anti-Hermitian, as \( \tilde{\mu} (\omega_m) = \tilde{\mu}_H (\omega_m) + \tilde{\mu}_{AH} (\omega_m) \) [46], to obtain:

\[
U_m = \frac{1}{2} \sum_{n=-\infty}^{+\infty} \sum_{m=-\infty}^{+\infty} \tilde{H}_n^* (\mathbf{r}) \cdot \left( i \omega_m \tilde{\mu}_H (\omega_m) - i \omega_m \tilde{\mu}_{AH} (\omega_m) \right) \cdot \tilde{H}_m (\mathbf{r}) \exp \left( i (\omega_m - \omega_n) t \right)
\]

\[
+ \frac{1}{2} \sum_{n=-\infty}^{+\infty} \sum_{m=-\infty}^{+\infty} \tilde{H}_n^* (\mathbf{r}) \cdot \left( i \omega_m \tilde{\mu}_{AH} (\omega_m) + i \omega_m \tilde{\mu}_H^* (\omega_m) \right) \cdot \tilde{H}_m (\mathbf{r}) \exp \left( i (\omega_m - \omega_n) t \right) \tag{A7}
\]

The same can be done for the electric energy density \( (U_e) \). In order to account for the real valued-ness of these quantities, one can break the summations into four parts as \( \sum_{m=-\infty}^{+\infty} \rightarrow \sum_{m=0}^{0} + \sum_{m=0}^{+\infty} \) and \( \sum_{n=-\infty}^{+\infty} \rightarrow \sum_{n=0}^{0} + \sum_{n=+\infty}^{+\infty} \) and change the indices in the first parts as \( m \rightarrow -m \) and \( n \rightarrow -n \). Carrying out the same procedure for all of the terms in equation (A5), one obtains the following term for the conservation of energy:

\[
\text{Re} \left\{ \tilde{E}_0^* (\mathbf{r}) \cdot \tilde{J}_0 (\mathbf{r}) + \mathbf{\nabla} \cdot \left( \tilde{E}_0 (\mathbf{r}) \times \tilde{H}_0^* (\mathbf{r}) \right) \right\}
\]

\[
+ \sum_{n=1}^{+\infty} \sum_{m=-n}^{+\infty} \left( \tilde{E}_n^* (\mathbf{r}) \cdot \tilde{J}_{n+1} (\mathbf{r}) + \mathbf{\nabla} \cdot \left( \tilde{E}_n (\mathbf{r}) \times \tilde{H}_{n+1}^* (\mathbf{r}) \right) \right)
\]
This equation is the general equation of the conservation of the spontaneous power for the electromagnetic pulses, stating that spontaneous power can be divided into two lower frequency and higher frequency parts, with harmonic frequencies denoted by \( \omega_m = -\omega_m \) and \( \omega_h = \omega_m + \omega_n \), respectively. If the pulses are quasi monochromatic and the repetition rate is so slow that one can assume \( T \to +\infty \), in a similar approach to [38] one can expand the terms \( \omega_m \) with a Taylor series and attribute the corresponding terms to the derivatives in the material dispersion.

A well-defined power spectrum can be obtained by averaging the spontaneous power over the fundamental period as

\[
\bar{P}_\text{sp} (\omega_n) = \frac{1}{T} \int_{-T/2}^{T/2} P_\text{sp} (\omega, t) dt.
\]

With such an assumption and using the identity

\[
\int_{-T/2}^{T/2} \exp \left( i (\omega_n - \omega_{m}) t \right) dt = 0,
\]

also neglecting the zero-order term, we obtain the following identity for the average power spectrum:

\[
\sum_{n=1}^{+\infty} \omega_n \left\{ \frac{1}{\omega_n} \iiint_s \text{Re} \left\{ \bar{E}_n (\bar{r}) \times \bar{H}_n^* (\bar{r}) \right\} \cdot ds \right\} - \iiint_v \bar{E}_n (\bar{r}) \cdot \bar{H}_n (\bar{r}) \cdot \text{Im} \left\{ \hat{\mu}_\text{AH} (\omega_n) \right\} \cdot \bar{H}_n^* (\bar{r}) \cdot dv \right\} = 0 \quad (A9)
\]

Note that the integration over the volume of interaction has been done explicitly. Since the summation should be zero for all of the frequency components, an obvious answer to equation (A9) is:
\[ -\frac{1}{\omega_n} \iiint_v \text{Re} \left\{ \vec{E}(\vec{r}, \omega_n) \cdot \vec{J}^*(\vec{r}, \omega_n) \right\} \, dv = \frac{1}{\omega_n} \iiint_s \text{Re} \left\{ \vec{E}(\vec{r}, \omega_n) \times \vec{H}^*(\vec{r}, \omega_n) \right\} \cdot ds \]
\[ - \iiint_v \vec{E}(\vec{r}, \omega_n) \cdot \text{Im} \left\{ \hat{e}_{\text{MH}}(\omega_n) \right\} \cdot \vec{E}^*(\vec{r}, \omega_n) \, dv \]
\[ - \iiint_v \vec{H}(\vec{r}, \omega_n) \cdot \text{Im} \left\{ \hat{\mu}_{\text{MH}}(\omega_n) \right\} \cdot \vec{H}^*(\vec{r}, \omega_n) \, dv \]

(A10)

It is interesting to note that the effective energy densities being stored at the volume of the interaction, in the forms of capacitive and inductive energies, do not contribute to the power spectrum averaged over the fundamental period, but only the dissipative terms appear in equation (A10). However, it should be stated that the detected power is in general related to the impulse response of the detector, and the averaging over the spontaneous power should be respectively carried out considering the acquisition rate of the detectors.

It is easy to demonstrate that the first term in equation (A10) is equivalent to the electron energy-loss spectra \( \Gamma_{\text{EELS}}(\omega) \) given in [37] within the undepleted pump approximation, only differing in an \( \hbar \) coefficient. Considering the photon generation spectra given by:

\[ \Gamma_{\text{PG}}(\omega) = \frac{1}{\hbar \omega} \iiint_s \text{Re} \left\{ \vec{E}(\vec{r}, \omega) \times \vec{H}^*(\vec{r}, \omega) \right\} \cdot ds \quad (A11) \]

and the EELS spectra given by:

\[ \Gamma_{\text{EELS}}(\omega) = -\frac{1}{\hbar \omega} \iiint_v \text{Re} \left\{ \vec{E}(\vec{r}, \omega) \cdot \vec{J}^*(\vec{r}, \omega) \right\} \, dv \quad (A12) \]

equation (A10) is simplified into:

\[ \Gamma_{\text{EELS}}(\omega) = \Gamma_{\text{PG}}(\omega) \]
\[ - \frac{1}{\hbar} \iiint_v \vec{H}(\vec{r}, \omega) \cdot \text{Im} \left\{ \hat{e}_{\text{MH}}(\omega) \right\} \cdot \vec{H}^*(\vec{r}, \omega) \, dv \]
\[ - \frac{1}{\hbar} \iiint_v \vec{E}(\vec{r}, \omega) \cdot \text{Im} \left\{ \hat{\mu}_{\text{MH}}(\omega) \right\} \cdot \vec{E}^*(\vec{r}, \omega) \, dv \quad (A13) \]

By assuming that the acquisition rate of the detector is equivalent to the emission rate of the source or is locked to that, the PG and EELS spectra only differ in the dissipative losses occurring in the system. It should also be mentioned that the results derived here are in very good agreement with the results reported by S. Glasgow et al [47]; however to derive equation (A13), the material properties are not restricted here to symmetric tensorial behaviors, and also the Kramers–Kronig relation is not used here. The only necessary assumption is the real valuedness of the field components and the tensors, to derive equation (A5).

References

[1] Goryashko V A and Ziemann V 2013 Self-amplified coherent spontaneous emission in a free electron laser with ‘quiet’ bunches Phys. Rev. Special Top. 16 030702
[2] McNeil B W J and Thompson N R 2010 X-ray free-electron lasers Nat. Photonics 4 814–21
[3] Corde S, Ta Phuoc K, Lambert G, Fitour R, Malka V, Rousse A, Beck A and Lefebvre E 2013 Femtosecond x rays from laser-plasma accelerators Rev. Mod. Phys. 85 1–48
[4] Killoran J H, Hacker F L and Walsh J E 1994 A millimeter wavelength radiation source using a dual grating resonator IEEE T. Plasma Sci. 22 530–5
[5] Levush B, Antonsen T M, Bromborsky A, Lou W R and Carmel Y 1992 Theory of relativistic backward-wave oscillators with end reflections Ieee T. Plasma Sci. 20 263–80
[6] Griguras I et al 2012 Ultrafast x-ray pulse characterization at free-electron lasers Nat. Photonics 6 852–7
[7] Service R F 2002 Battle to become the next-generation x-ray source Science 298 1356–8
[8] Pfeifer T, Spielmann C and Gerber G 2006 Femtosecond x-ray science Rep. Prog. Phys. 69 443–505
[9] Doucas G, Mulvey J H, Omori M, Walsh J and Kimmitt M F 1992 1st observation of smith-purcell radiation from relativistic electrons Phys. Rev. Lett. 69 1761–4
[10] Urata J, Goldstein M, Kimmitt M F, Naumov A, Platt C and Walsh J E 1998 Superradiant smith-purcell emission Phys. Rev. Lett. 80 516–9
[11] de Abajo F J G 2000 Smith-purcell radiation emission in aligned nanoparticles Phys. Rev. E 61 5743–52
[12] Adamo G, MacDonald K F, Zheludev N I, Fu Y H, Wang C M, Tsai D P and de Abajo F J G 2009 Tunable, nanoscale free-electron source of photons and plasmons Conf. on Lasers and Electro-Optics and Quantum Electronics and Laser Science Conf. (Cleo/Qels) vol 1–5 pp 2479–80
[13] Adamo G, MacDonald K F, Fu Y H, Wang C M, Tsai D P, de Abajo F J G and Zheludev N I 2009 Light well: a tunable free-electron light source on a chip Phys. Rev. Lett. 103 1113901
[14] Adamo G, Ou J Y, So J K, Jenkins S D, de Angelis F, MacDonald K F, Di Fabrizio E, Ruostekoski J and Zheludev N I 2012 Electron-beam-driven collective-mode metamaterial light source Phys. Rev. Lett. 109 217401
[15] Lounis B and Orrit M 2005 Single-photon sources Rep. Prog. Phys. 68 1129–79
[16] Raimond J M, Brune M and Haroche S 2001 Colloquium: manipulating quantum entanglement with atoms and photons in a cavity Rev. Mod. Phys. 73 565–82
[17] Yuan Z L, Kardynal B E, Stevenson R M, Shields A J, Lobo C J, Cooper K, Beattie N S, Ritchie D A and Pepper M 2002 Electrically driven single-photon source Science 295 102–9
[18] Akimov A V, Mukherjee A, Yu C L, Chang D E, Zibrov A S, Hammer P R, Park H and Lukin M D 2007 Generation of single optical plasmons in metallic nanowires coupled to quantum dots Nature 450 402–6
[19] Curto A G, Volpe G, Taminiau T H, Kreuzer M P, Quindant R and van Hulst N F 2010 Unidirectional emission of a quantum dot coupled to a nanoantenna Science 329 930–3
[20] Gazzano O, de Vasconcellos S M, Arnold C, Nowak A, Galopin E, Sagnes I, Lanco L, Lemaitre A and Senellart P 2013 Bright solid-state sources of indistinguishable single photons Nat. Commun. 4 1425
[21] de Martini F, Di Giuseppe G and Marrocco M 1996 Single-mode generation of quantum photon states by excited single molecules in a microcavity trap Phys. Rev. Lett. 76 900–3
[22] Lounis B and Moerner W E 2000 Single photons on demand from a single molecule at room temperature Nature 407 491–3
[23] Kursruefer C, Mayer S, Zarda P and Weinfurter H 2000 Stable solid-state source of single photons Phys. Rev. Lett. 85 290–3
[24] Geiselmamn M, Juan M L, Renger J, Say J M, Brown L J, de Abajo F J G, Koppens F and Quindant R 2013 Three-dimensional optical manipulation of a single electron spin Nat. Nanotechnol. 8 175–9
[25] Koyama K, Yoshita M, Baba M, Suemoto T and Akiyama H 1999 High collection efficiency in fluorescence microscopy with a solid immersion lens Appl. Phys. Lett. 75 1667–9
[26] Lukosz W 1979 Light-emission by magnetic and electric dipoles close to a plane dielectric interface.3. radiation-patterns of dipoles with arbitrary orientation J. Opt. Soc. Am. 69 1495–503
[27] Lee K G, Chen X W, Eghlidi H, Kukura P, Lettow R, Renn A, Sandoghdar V and Gotzinger S 2011 A planar dielectric antenna for directional single-photon emission and near-unity collection efficiency Nat. Photonics 5 166–9
[28] Lee K G, Eghlidi H, Chen X W, Renn A, Gotzinger S and Sandoghdar V 2012 Spontaneous emission enhancement of a single molecule by a double-sphere nanoantenna across an interface Opt. Express 20 23331–8
[29] Schietinger S, Barth M, Alchele T and Benson O 2009 Plasmon-enhanced single photon emission from a nanoassembled metal-diamond hybrid structure at room temperature Nano Lett. 9 1694–8
[30] Taminiau T H, Stefani F D, Segerink F B and Van Hulst N F 2008 Optical antennas direct single-molecule emission Nat. Photonics 2 234–7
[31] Englund D, Shields B, Rivoire K, Hatami F, Vuckovic J, Park H and Lukin M D 2010 Deterministic coupling of a single nitrogen vacancy center to a photonic crystal cavity Nano Lett. 10 3922–6
[32] Bendana X, Polman A and de Abajo F J G 2011 Single-photon generation by electron beams Nano Lett. 11 5099–103
[33] Tizei L H G and Kociak M 2013 Spatially resolved quantum nano-optics of single photons using an electron microscope Phys. Rev. Lett. 110 153604
[34] Baum P 2013 On the physics of ultrashort single-electron pulses for time-resolved microscopy and diffraction Chem. Phys. 423 55–61
[35] Gliserin A, Apolonski A, Krausz F and Baum P 2012 Compression of single-electron pulses with a microwave cavity New J. Phys. 14 073055
[36] Talebi N, Sigle W, Vogelgesang R and Van Aken P 2013 Numerical simulations of interference effects in photon-assisted electron energy-loss spectroscopy New J. Phys. 15 053013
[37] de Abajo F J G 2010 Optical excitations in electron microscopy Rev. Mod. Phys. 82 209–75
[38] Jackson J D 1998 Classical Electrodynamics (USA: John Wiley and Sons, Inc.)
[39] Harrington R F 1961 Time-Harmonic Electromagnetic Fields (New York: McGraw-Hill, Inc.)
[40] Garcia Molina R, Grasmarti A, Howie A and Ritchie R H 1985 Retardation effects in the interaction of charged-particle beams with bounded condensed media J. Phys. C Solid State 18 5335–45
[41] Etchegoin P G, Le Ru E C and Meyer M 2007 Erratum: ‘An analytic model for the optical properties of gold’ [J. Chem. Phys. 125 164705 (2006)] J. Chem. Phys. 127 189901
[42] Etchegoin P G, Le Ru E C and Meyer M 2006 An analytic model for the optical properties of gold J. Chem. Phys. 125 164705
[43] Talebi N, Sigle W, Vogelgesang R, Koch C T, Fernandez-Lopez C, Liz-Marzán L M, Ogut B, Rohm M and van Aken P A 2012 Breaking the mode degeneracy of surface plasmon resonances in a triangular system Langmuir 28 8867–73
[44] Ogut B, Talebi N, Vogelgesang R, Sigle W and van Aken P A 2012 Toroidal plasmonic eigenmodes in oligomer nanocavities for the visible Nano Lett. 12 5239–44
[45] Esslinger M, Khunsin W, Talebi N, Wei T, Dorfmüller J, Vogelgesang R and Kern K 2013 Phase engineering of subwavelength unidirectional plasmon launchers Adv. Opt. Mater. 1 434–7
[46] Esslinger M and Vogelgesang R 2012 Reciprocity theory of apertureless scanning near-field optical microscopy with point-dipole probes Acs Nano 6 8173–82
[47] Glasgow S, Ware M and Peatross J 2001 Poynting’s theorem and luminal total energy transport in passive dielectric media Phys. Rev. E 64 046610