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Emission spectrum of a mesoscopic ring driven by fast unipolar pulses

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Abstract

We perform a theoretical study of the emission properties of a thin ballistic, mesoscopic ring driven by a quasi-periodic train of electromagnetic half-cycle pulses (HCPs). It is shown that an applied sequence of pulses induces a non-equilibrium, time-dependent charge polarization in the ring. The resulting charge oscillations generate an emission spectrum that can be modified by designing appropriately the strength and the shape of the HCP sequence. The HCP is fast in that, the relaxation of the excited ring to its equilibrium state occurs after the pulses have passed by. Hence the study of the emission spectrum offers a novel possibility to investigate relaxation processes in mesoscopic systems in absence of external perturbations.

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1. Introduction

Due to their importance for technological applications as well as for fundamental research, mesoscopic systems are continuing to attract considerable attention [1–4]. Of particular interest for this work are mesoscopic systems with ring-confining geometries. At low temperatures, the phase coherence length of electrons in a mesoscopic ring (MR) can be considerably large compared to the size of the ring and interference effects become important. Interference effects are manifested in various phenomena. E.g., when a MR is threaded by a magnetic field, the thermodynamic properties of the system become periodic functions of the magnetic flux. The flux dependence of the free energy leads then to the existence of the so-called persistent currents [4–9,12] and to the Aharonov–Bohm conductance oscillations [1].

Charge transfer phenomena and the polarizability of MRs have been studied previously both theoreti-

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cally [10] and experimentally [11]. In these studies, however, MRs are considered in the presence of a time-independent flux or dc electric fields, i.e., the system behaviour is investigated in the presence of external perturbations. In addition, the rise-up and switch-off times of these external fields occur usually adiabatically on the scale of the relaxation time $\tau_{\rm rel}$. Thus, a time-resolved study of the unperturbed system with a time resolution better than $\tau_{\rm rel}$ is not feasible (this includes in particular the dynamical behaviour of the relaxation itself). In the present Letter we address the situation in which a MR is periodically excited within times shorter than $\tau_{\rm rel}$ and relaxes to its equilibrium state in a field-free manner after each perturbation. Specifically, we investigate the charge dynamics and the emission spectrum of a thin (one subband) ballistic MR subjected to a train of linearly polarized electromagnetic half-cycle pulses (HCPs). An HCP is a strongly asymmetric mono-cycle pulse consisting of a very short, strong half-cycle (we refer to this part as an HCP), followed by a second half-cycle of an opposite polarity (the tail of the HCP). This second part of the pulse can be attenuated and stretched substantially (in time) by means of optical gating techniques [13]. Since the HCP tail is very weak and very long (compared to the relaxation time of MRs) it hardly influences the electron dynamics. Nowadays, subpicosecond HCPs are available that last for picoseconds and have a peak field up to several hundreds of kV/cm [14]. Furthermore, trains of HCPs are experimentally feasible [15–17].

The highly asymmetric HCPs and symmetric (or nearly symmetric) fields (e.g., continuous wave (CW) lasers and laser pulses) differ qualitatively in the way they couple to electronic systems. A key difference is that unlike CW lasers or laser pulses, an HCP can deliver a non-zero momentum transfer (or a *kick*) to the system [13,16–18]. It is precisely this peculiar property of HCPs that allows for inducing postpulse (and, therefore, field-free) time-dependent charge oscillations in MRs, as shown below. When a MR is illuminated by a single HCP the charge oscillations decay on a time scale of the order of the relaxation time $\tau_{\rm rel}$. However, the charge oscillations can be sustained for a longer time if a train of HCPs is applied. The induced charge oscillations correspond to a characteristic emission, meaning that the driven MR can serve as a source of electromagnetic radiation. As shown here, the emission characteristics of the MR can be controlled (to a certain extent) by appropriately designing the sequence of HCPs.

2. Theory

We consider an isolated 1D MR at low temperatures ($T \approx 0$ K). A periodic train of HCPs linearly polarized along the x direction is applied at $t = 0$. We are particularly interested in the case where the period T of the pulse sequence is much longer than the (relaxation) time $\tau_{\rm rel}$ which is needed for the excited MR to relax to the equilibrium state after the action of each pulse. In such a case the task is reduced to considering the evolution of the system under the action of a single pulse. The system evolution under the influence of the HCP train can be retrieved from the behaviour under the action of a single HCP. In addition, the duration Δt of an HCP is assumed to be much shorter than the ballistic time τ_F a particle at the Fermi level requires for completing one turn around the ring. This condition is experimentally feasible: for typical ballistic rings τ_F is of the order of several tens of picoseconds (here we consider a ring with $\tau_F \approx 32$ ps as the experimental ring used in Ref. [9]), while HCPs as short as 1 ps can be generated with contemporary techniques [14]. For $\Delta t \ll \tau_F$ the interaction of the system with the HCPs can be treated justifiably within the impulsive approximation. As specified below, the MRs we are considering are such that in the equilibrium state the charge carrier in the ring can be treated independently [8,9]. Thus, the fundamental quantity from which the behaviour of the complete system can be deduced is the single-particle, time-dependent wave function Ψ. The function Ψ is deduced from the solution of the time-dependent Schrödinger equation in the presence of a single HCP, i.e.,

$$i\hbar\frac{\partial\Psi}{\partial t} = \left[-\frac{\hbar^2}{2m^*\rho_0^2}\frac{\partial^2}{\partial\theta^2} - q\rho_0 p\cos\theta\,\delta(t)\right]\Psi. \qquad (1)$$

Here ρ_0 is the radius of the ring and $\delta(t)$ denotes the Dirac δ function. The polar angle θ specifies the angular position of the charge carrier (with charge q) with respect to the x axis and m^* is the effective mass. The pulse is applied at $t = 0$ and transfers a momentum

denoted by p to the system [14,16–18]. The momentum p is determined by the area of the actual pulse, i.e., by the time-integral $\int_0^{\Delta t} F(t)\,dt$ over the electric field amplitude $F(t)$ of the pulse. Regarding (1), it is worth remarking that although the actual HCP has a finite duration its approximation by a Dirac delta function (this is equivalent to the impulsive approximation) is expected to describe appropriately the dynamics of the system if the pulse duration is much shorter than the characteristic time (in the present context τ_F) of the field-free system. In fact, it has been numerically demonstrated in a number of works (see, for example, Ref. [18]) that under the mentioned condition the impulsive approximation leads indeed to results that hardly differ from the numerical results obtained when considering the finite duration of the HCP.

It is worth noting that the 1D model is expected to be a good approximation for thin MRs with a width $d \ll \rho_0$. In such a case the radial motion is much faster than the angular motion and, hence, the radial channels become adiabatically decoupled from the angular motion. If, in addition, $d < \lambda_F$ (with λ_F representing the Fermi wavelength) then the MR becomes a single-channel 1D ring. A typical value for the Fermi wavelength in ballistic, thin GaAs–AlGaAs MRs is $\lambda_F = 42$ nm (see, for example, Ref. [9]).

The time evolution of a particle residing initially (at $t < 0$) in the electronic state characterized by the rotational quantum number m_0 is described by the solution $\Psi_{m_0}(\theta, t)$ of Eq. (1) at $t > 0$. Expanding $\Psi_{m_0}(\theta, t)$ in terms of the stationary eigenstates leads to

$$\Psi_{m_0}(\theta, t) = \frac{1}{\sqrt{2\pi}} \sum_{m=-\infty}^{\infty} C_m(t, m_0) e^{im\theta} e^{-i\frac{E_m t}{\hbar}}, \quad (2)$$

where

$$E_m = \frac{\hbar^2 m^2}{2m^* \rho_0^2}, \quad m = 0, \pm 1, \pm 2, \ldots \quad (3)$$

are the eigenenergies of the unperturbed states. The expansion coefficients $C_m(t, m_0)$ have the form

$$C_m(t, m_0) = \begin{cases} \delta_{m,m_0} & \text{for } t \leqslant 0, \\ i^{m_0 - m} J_{m-m_0}(\alpha) & \text{for } t > 0, \end{cases} \quad (4)$$

where $J_l(x)$ are the Bessel functions and $\alpha = q\rho_0 p/\hbar$. On the other hand, the energy corresponding to a particle initially in the m_0th state is derived to be

$$E_{m_0}(t) = \begin{cases} \frac{\hbar^2 m_0^2}{2m^* \rho_0^2} & \text{for } t \leqslant 0, \\ \frac{\hbar^2}{2m^* \rho_0^2}\left(m_0^2 + \frac{\alpha^2}{2}\right) & \text{for } t > 0. \end{cases} \quad (5)$$

Thus, the HCP shifts the unperturbed energy spectrum by an amount that scales quadratically with the strength of the pulse and does not depend on the size of the ring.

The polarization (along the x axis) of the MR induced by the HCP is given by

$$\mu(t) = \sum_{m_0, \sigma} f(m_0, \sigma, N, t) \mu_{m_0}(t). \quad (6)$$

The spin of the particle is referred to by σ and N is the number of particles (which is assumed to be a constant since the ring is isolated). In Eq. (6) $f(m_0, \sigma, N, t)$ stands for the non-equilibrium distribution function. The quantity μ_{m_0}, defined as

$$\mu_{m_0}(t) = q\rho_0 \int_0^{2\pi} \left|\Psi_{m_0}(\theta, t)\right|^2 \cos\theta\, d\theta \quad (7)$$

is the dipole moment along the x axis associated with a particle that resided initially in the m_0th stationary state.

As shown in Ref. [8], the electron–electron interaction plays a minor role in determining the ground state properties of the MRs considered in this Letter. When the HCP is applied to the MR, the system is promoted into an excited state and starts to relax after the HCP has passed by. The relaxation of the system can occur through various pathways, e.g., electron–phonon scattering, simultaneous scattering by impurities and phonons, electron–electron scattering,[1] etc. The detailed study of such mechanisms, although interesting, is out of the scope of the present Letter. Here the relaxation processes are introduced phenomenologically by means of a single (averaged) quantity, the relaxation

[1] We note that unlike in the equilibrium state in which the charge is distributed uniformly along the ring and the electron–electron scattering conserves the momentum in the two-particle collisions, in the excited state the external field induces a polarization that breaks the rotational symmetry and leads to the violation of the momentum conservation. A similar situation occurs in "dirty" metal rings in which the presence of disorder is responsible for breaking the rotational symmetry.

time $\tau_{\rm rel}$. The non-equilibrium distribution function $f(m_0,\sigma,N,t)$ in Eq. (6) is evaluated within the relaxation time approximation by solving the Boltzmann equation:

$$\frac{\partial f(m_0,t)}{\partial t} = -\frac{f(m_0,t)-n_F(m_0)}{\tau_{\rm rel}}, \tag{8}$$

where $\tau_{\rm rel}$ represents the relaxation time and

$$n_F(m_0) = \left[1+\exp\left(\frac{E_{m_0}(t\leqslant 0)-\eta_0}{k_BT}\right)\right]^{-1} \tag{9}$$

denotes the Fermi–Dirac distribution function corresponding to the equilibrium. Note that the explicit dependence on σ and N has been omitted for brevity. In the equation above T, k_B, and η_0 represent the temperature, the Boltzmann constant and the chemical potential (for $t\leqslant 0$), respectively. Eq. (8) has to be complemented with the boundary condition specifying the value of the distribution function right after the application of the pulse,

$$f\left(m_0,0^+\right) = n_F^{(1)}(m_0) = \left[1+\exp\left(\frac{E_{m_0}(t>0)-\eta_1}{k_BT}\right)\right]^{-1}. \tag{10}$$

In the present case of an isolated MR, the chemical potentials η_0 and η_1 have to be calculated, in general, by requiring the number of particles N in the ring to be a constant. In the particular case of a zero temperature, only the lowest-lying states are occupied. In such a case, considering that the action of the HCP on the energy spectrum consists only in shifting the energy levels (see Eq. (5)), the sum in Eq. (6) can be performed analytically. Hereby the filling of the levels depending on whether the N charge carriers are spinless or spin $\frac{1}{2}$ particles plays a crucial role. For the case of spinless particles [$\mu(\sigma=0)=\mu^0$] we obtain

$$\mu^0 = \begin{cases} \mu_o(N,t) & \text{for } N \text{ odd},\\ \mu_e(N,t) & \text{for } N \text{ even}, \end{cases} \tag{11}$$

where

$$\mu_o(N,t) = -\alpha q\rho_0\,\Theta(t)\left[J_0(\Omega)+J_2(\Omega)\right] \times \sin\left[\frac{2\pi Nt}{\tau_p}\right]e^{-t/\tau_{\rm rel}}, \tag{12}$$

$$\mu_e(N,t) = -q\alpha\rho_0\,\Theta(t)\left[J_0(\Omega)+J_2(\Omega)\right] \times \sin\left[\frac{2\pi Nt}{\tau_p}\right]\cos\left[\frac{2\pi t}{\tau_p}\right]e^{-t/\tau_{\rm rel}}. \tag{13}$$

In Eqs. (12) and (13) $\Theta(t)$ denotes the Heaviside step function, $\alpha = q\rho_0 p/\hbar$, $\tau_p = 4\pi m^*\rho_0^2/\hbar$ and $\Omega = \alpha\sqrt{2-2\cos[4\pi t/\tau_p]}$.

By following a procedure similar to that discussed in Ref. [7] the total HCP-induced dipole moment $\mu^\sigma(t)$ corresponding to the case of spin $\frac{1}{2}$ particles can be expressed in terms of the dipole moments for the spinless case (Eqs. (12) and (13))

$$\mu^\sigma = \begin{cases} 2\mu_e\left(\frac{N}{2},t\right) & \text{if } N=0 \pmod 4,\\ \mu_o\left(\frac{N+1}{2},t\right)+\mu_e\left(\frac{N-1}{2},t\right) & \text{if } N=1 \pmod 4,\\ 2\mu_o\left(\frac{N}{2},t\right) & \text{if } N=2 \pmod 4,\\ \mu_e\left(\frac{N+1}{2},t\right)+\mu_o\left(\frac{N-1}{2},t\right) & \text{if } N=3 \pmod 4. \end{cases} \tag{14}$$

We note that for time scales in the domain $t\ll\tau_p$ the dependence of the polarization on the parity of the number of particles N becomes irrelevant (compare Eqs. (12) and (13)). In such a situation one can easily find from Eqs. (12)–(14) that for a large value of N [$(N\pm 1)/N\approx 1$], all the four cases in Eq. (14) can be approximated as follows

$$\mu^\sigma(N,t) \approx -2\alpha q\rho_0\,\Theta(t)\left[J_0(\Omega)+J_2(\Omega)\right] \times \sin\left[\frac{\pi Nt}{\tau_p}\right]e^{-t/\tau_{\rm rel}}. \tag{15}$$

The polarization $\mu_k(t)$ (for the case of spin $\frac{1}{2}$ particles) of the MR due to a periodic sequence of k HCPs with period $T\gg\tau_{\rm rel}$ can be obtained from the polarization induced by a single HCP [Eq. (15)] through the following relation

$$\mu_k(t) = \sum_{j=0}^{k-1}\mu^\sigma(t-kT), \tag{16}$$

where the dependence on N has been omitted for brevity.

The emission spectrum $I(\omega)$ produced by the charge oscillations in the MR is given by

$$I(\omega)\sim\left|\mu_k(\omega)\right|^2, \tag{17}$$

where $\mu_k(\omega)$ is obtained by Fourier transforming $\mu_k(t)$ to the frequency (ω) domain. In the limit of weak fields ($\alpha\ll 1/2$) we have $\Omega\ll 1$. The Bessel

functions in Eq. (15) can then be approximated as [19]

$$J_\nu(\Omega) \approx \frac{1}{\Gamma(\nu+1)}\left(\frac{\Omega}{2}\right)^\nu, \quad \nu \neq -1,-2,-3,\ldots. \tag{18}$$

Consequently, from Eqs. (15)–(18) the following analytical expression is deduced

$$I(\omega) \sim Y(\omega)\left[\frac{\sin\left(\frac{k\omega T}{2}\right)}{\sin\left(\frac{\omega T}{2}\right)}\right]^2, \tag{19}$$

where the function $Y(\omega)$ is given by

$$Y(\omega) = \frac{(2\pi\alpha q\rho_0 N)^2}{\left(\frac{\tau_p}{\tau_{\rm rel}^2} + \frac{\pi^2N^2}{\tau_p} - \omega^2\tau_p\right)^2 + 4\omega^2\frac{\tau_p^2}{\tau_{\rm rel}^2}}. \tag{20}$$

One can see from Eqs. (19) and (20) that at the integer harmonics ($\omega = n\omega_0, \omega_0 = 2\pi/T, n = 0, 1, 2, \ldots$) the emission spectrum has peaks whose amplitudes are modulated by the modulation function $k^2Y(\omega)$. The modulation function has a maximum

$$k^2Y_{\max} = (k\alpha q\rho_0\tau_{\rm rel})^2 \tag{21}$$

at the frequency

$$\frac{\omega_{\max}}{\omega_0} = \sqrt{\left(\frac{NT}{2\tau_p}\right)^2 - \left(\frac{T}{2\pi\tau_{\rm rel}}\right)^2}. \tag{22}$$

Note, however, that as $\omega_{\max}$ must be a real number, the maximum indicated in Eq. (21) exists only if $\tau_{\rm rel} \geqslant \tau_p/(\pi N)$.

3. Results

Explicit calculations were performed for a ballistic GaAs–AlGaAs ring similar to that used in the experiment reported in Ref. [9] with the following parameters: $\rho_0 = 1.35$ μm, the electron effective mass $m^* = 0.067m_e$, and $N = 1400$. For such a system $\tau_p \approx 13.26$ ns. The time domain of interest here is $t \ll \tau_p$, where Eq. (15) and the further relations derived from it hold. A quasiperiodic train of $k = 10$ sine-square HCPs with a period $T = 100$ ps are used. The HCPs were taken as short as 1 ps and with peak field as weak as $F = 1$ V/cm (to ensure that $\alpha \ll 1/2$ holds). A zero temperature was considered in all calculations.

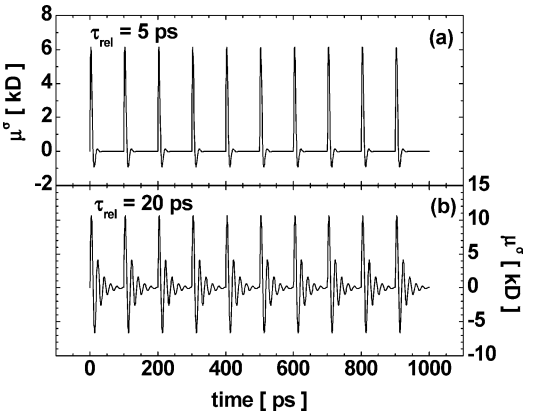

Fig. 1. Time dependence of the polarization μ^σ (in units of 10^3 Debyes) for different values of $\tau_{\rm rel}$. The results correspond to a train of $k = 10$ HCPs with a period of $T = 100$ ps and a peak field of $F = 1$ V/cm.

The time dependence of the induced polarization μ^σ is shown in Fig. 1(a) and (b) for $\tau_{\rm rel} = 5$ ps and $\tau_{\rm rel} = 20$ ps, respectively. It is clear from the comparison of these figures that as the relaxation time increases the carriers in the ring are able to oscillate stronger and therefore the amplitude and the number of oscillations of the induced polarization increases with $\tau_{\rm rel}$. It is worth mentioning that although our study here is limited to the case of a quasiperiodic train of unidirectional HCPs, it is also experimentally feasible to generate trains of bidirectional HCPs [17] as well as to vary the time delay between consecutive pulses. Therefore, it is possible to control the ring polarization on a picosecond time scale by appropriately designing the sequence of HCPs.

The emission spectra corresponding to the polarization oscillations shown in Fig. 1(a) and (b) are displayed in Fig. 2(a) and (b), respectively. For the sake of comparison $I(\omega)$ has been expressed in the same arbitrary units in both Fig. 2(a) and (b). As anticipated in the previous section, the emission spectrum is composed of peaks at the integer harmonics ($\omega = n\omega_0, n = 0, 1, 2, \ldots$) with amplitudes that are modulated by the modulation function $k^2Y(\omega)$. This situation is in contrast to the harmonic spectra obtained from a single atom in a linearly polarized laser field and from ringlike molecules interacting with circularly polarized fields. In the former case only linearly polarized odd harmonics are generated [20] and in the latest only harmonics of the order $n = gK \pm 1$ (K is the or-

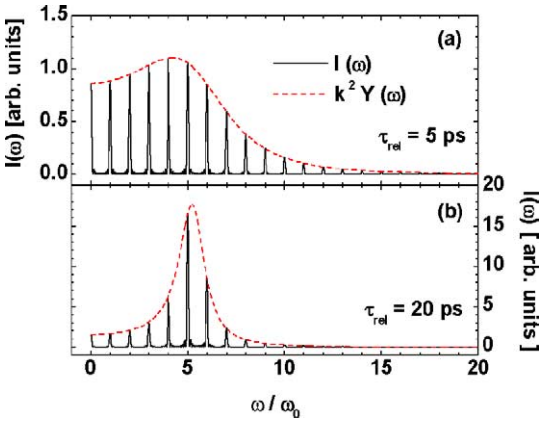

Fig. 2. Emission spectrum (solid lines) and modulation function (dashed lines) for different values of τ_{rel}. The results correspond to a train of $k = 10$ HCPs with a period of $T = 100$ ps ($\omega_0 = 2\pi/T \approx 6.28 \times 10^{10}$ Hz) and a peak field of $F = 1$ V/cm.

der of the discrete rotational symmetry of the molecule and $g = 1, 2, 3, \ldots$) can be emitted [21]. In these cases the suppression of determined harmonics is due to selection rules imposed by the dynamical symmetry of the corresponding Hamiltonian. For the system studied here the highly asymmetric character of the HCPs has the consequence that the Hamiltonian does not have any dynamical symmetry and therefore no special selection rules exist.

It is quite apparent that the intensity of the strongest emission lines increases with the relaxation time (compare Fig. 2(a) and (b)). This behaviour is comprehensible from Eq. (21), for the maximum of the modulation function depends quadratically on τ_{rel}. Furthermore, as the modulation function depends essentially on the ring parameters, by choosing appropriate values for ρ_0 and N one can *filter* (within a certain accuracy) some specific harmonics from the emission spectrum, e.g., one can choose the system parameters such that $\omega_{\max}/\omega_0$ (see Eq. (22)) coincides with the order of the harmonic one wishes to highlight. We stress that the emission of concern here is mainly due to the induced charge oscillations and not to the relaxation processes. Due to the single-channel (1D) condition the quantum transitions between different states are more likely to occur through the gain or loss of angular momentum and, therefore, the relaxation of the system is expected to occur through phonon (and not photon) emission. For the same reason the effects of the emitted photons on the relaxation processes are not expected to be important and, consequently, are not included in our model.

It is worth noting that the charge polarization effects occur mainly between consecutive pulses and, therefore, in a field-free environment. This fact offers a unique opportunity for studying relaxation processes in absence of external perturbations. In fact, the results depicted in Fig. 2 suggest a way of experimentally measuring the relaxation time τ_{rel}. By detecting the emission spectrum, the relaxation time can be found as the value of τ_{rel} in the modulation function $k^2 Y(\omega)$ that leads to the best correspondence with the emission peaks that are determined experimentally. In addition, we believe that it is, in principle, possible to infer from the emission spectrum whether there are several relaxation mechanisms and which of them are the most important. A definitive answer to this question, however, requires the detailed inclusion of the relaxation mechanisms (one can then "switch on" or off the different relaxation mechanisms and by comparing with the experimental results one can obtain information about the relaxation process). The model discussed here includes the relaxation at a phenomenological level and cannot provide such a detailed information.

The system studied here could be useful for harmonic generation as well as for designing artificially structured materials that do not occur naturally. In particular the creation of a planar array of isolated MRs[2] is expected to resonantly increase the emission intensity.[3] Furthermore, as shown above, the individual ring polarization strongly depends on its parameters (mainly size, and number of particles). Therefore, upon an appropriate design of the corresponding planar array (e.g., by varying the diameters of certain rings in the array) one can create a desired map of the charge polarization of the quasi two-dimensional structured material.

4. Conclusions

In summary, we showed that when a thin ballistic MR is subjected to a quasiperiodic train of HCPs

[2] Similar arrays of connected GaAs/GaAlAs MRs have been experimentally realized (see Ref. [12]).

[3] An experimental realization of an artificially structured material with an enhanced response has been reported in Ref. [22].

a non-equilibrium, time-dependent polarization is induced in the ring. The charge oscillations induced in this way can be engineered on the picosecond time scale by an appropriate design of the pulse sequence. The emission spectrum generated by the induced charge oscillations was studied for the case of weak pulses. The obtained results document the potential of the studied system for harmonic generations. A procedure was proposed for the experimental measurement of the time of relaxation of the excited states to the equilibrium state (in absence of external perturbations).

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