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Theoretical investigations of the switch-on behavior of semiconductor quantum dot based nanocavity laser devices are presented. From a microscopic treatment of the carrier-carrier and carrier-phonon interaction, we find a fast switch-on of the laser device that is enabled by ultrafast carrier dynamics and heavily damped relaxation oscillations. We show that the timescales of the dynamics within the continuum states and the quantum dot states are strongly coupled and investigate the time dependence of the non-equilibrium scattering rates in detail. © 2011 American Institute of Physics. [doi:10.1063/1.3651765]

Semiconductor lasers are central components of optical technologies. They are most prominently used in optical communications and optical data storage involving compact discs, digital versatile discs, and Blu-ray discs. Recently, semiconductor active regions coupled with nanoscale optical resonator schemes have opened an additional level of miniaturization and facilitated the implementation of photonic integrated circuits. The advent of photonic crystals allows to resonator schemes have opened an additional level of miniaturization and facilitated the implementation of photonic integrated circuits. The advent of photonic crystals allows to create high-quality cavities,1–4 which opens a multitude of possibilities for guiding and modifying the emission properties via Purcell enhancement of emission rates,4 and also allows to directly integrate nanocavities with optical waveguides. For such photonic circuits, nanolasers are an important ingredient. The need for ever-higher data transmission rates calls for an increase in the modulation speed of the underlying laser devices. A fast switch-on behavior is a key quantity towards the realization of this goal, as unfavorable switch-on characteristics significantly limit the maximum possible transmission rate.

In recent years, semiconductor quantum dots (QDs) have proven to be an interesting system, as they possess a discrete energy spectrum that can be engineered to a large extent. They are considered as gain material for next generation optoelectronic devices and for fundamental studies of light-matter interaction.5–7

For the switch-on and modulation properties of QD laser devices, carrier scattering is a key quantity. Previous investigations on carrier scattering 8 and its influence on the dynamical properties of QD lasers9–11 were based on quasi-equilibrium evaluations or rate equation analysis. In contrast, we present in this letter a full non-equilibrium analysis of the switch-on properties of QD-based nanolaser systems, including an analysis of carrier heating and non-equilibrium scattering rates.

We consider an ensemble of InGaAs QDs embedded in an optical nanocavity. Taken into account are the two lowest confined shells for electrons and holes, that are labeled s- and p-shell, respectively, as well as the continuum of wetting layer (WL) states. Within the cluster expansion scheme, we derive coupled equations of motion for quantities like the photon number in the cavity mode and the carrier population of the different QD and WL states, extending previous work that focused on the QD states and carrier-photon correlations.12 It should be noted that, contrary to other investigations9,10 we focus on nanocavity devices with a well separated mode spectrum and strong Purcell enhancement. The equations for the population of the state (QD or WL) and for the photon number in the mode are

\[
\frac{\hbar}{d} f^j_v = -2 \text{Re} \sum_q |g_{qv}|^2 \langle b^j_q v_i^c c_i \rangle
\]

\[
+ S^j_{v, \text{Corr}} + (1 - f^c - f^h) P_v,
\]

\[
\left( \frac{\hbar}{d} + 2 \kappa_q \right) n_q = +2 \text{Re} \sum_v |g_{qv}|^2 \langle b^j_q v_i^c c_v \rangle.
\]

Here, \( \kappa_q \) is the cavity loss rate, connected to the Q-factor, \( g_{qv} \) is the light-matter coupling constant, and \( P_v \) is the pump contribution for the state \( v \). Both \( n_q \) and \( f^c \) couple to the photon assisted polarization \( \langle b^j_q v_i^c c_v \rangle \)

\[
\left( \frac{\hbar}{d} + \kappa_q + \Gamma + i (\zeta^e + \zeta^h - \hbar \omega_q) \right)
\]

\[
\times \langle b^j_q v_i^c c_v \rangle = f^c f^h - (1 - f^c - f^h) n_q
\]

\[
+ \delta \langle b^j_q b^j_q c_i c_i \rangle - \delta \langle b^j_q b^j_q v_i^c v_i^c \rangle.
\]

While it has been shown that an accurate description of dephasing under lasing conditions requires a non-Markovian treatment of the carrier-carrier and carrier-phonon interaction mechanisms,13 we include its influence into an effective dephasing rate \( \Gamma \), which is determined by the carrier density dependent fit to the full non-Markovian calculation.14 A full microscopical inclusion of these effects is expected to only give minor changes as the emission dynamics is not very sensitive towards dephasing.

The higher order correlation functions, \( \delta \langle b^j_q b^j_q c_i c_i \rangle \) and \( \delta \langle b^j_q b^j_q v_i^c v_i^c \rangle \), obey their own equations of motion, see, e.g., Ref. 12 for details.

The rate \( S^j_{v, \text{Corr}} \) includes carrier-carrier and carrier-phonon scattering and is given by

\[
S^j_{v, \text{Corr}} = -f^j_v S^j_{v, \text{out}} + (1 - f^c - f^h) S^j_{v, \text{in}} - \frac{f^j_v - F^j_v}{t_{\text{phonon}}}.
\]
Here, the rates \( S_i^{\text{in, out}} \) for the carrier-carrier Coulomb interaction are evaluated in second order Born approximation and can be written as

\[
S_i^{\text{in}} = \frac{2}{\hbar} \sum_{\mu_1 \mu_2 \nu_1 \nu_2} \langle \rho_{\mu_2 \nu_2} \rangle \left( W_{\mu_2 \nu_2 \mu_1 \nu_1} - W_{\nu_2 \nu_1 \mu_1 \mu_2} \right) \times f_i^{\mu_1} (1 - f_i^{\mu_2}) \delta(\varepsilon_i^{\mu_1} - \varepsilon_i^{\mu_2})
\]

and

\[
S_i^{\text{out}} = \frac{2}{\hbar} \sum_{\mu_1 \mu_2 \nu_1 \nu_2} \langle \rho_{\mu_2 \nu_2} \rangle \left( W_{\mu_2 \nu_2 \mu_1 \nu_1} - W_{\nu_2 \nu_1 \mu_1 \mu_2} \right) \times (1 - f_i^{\mu_1}) f_i^{\mu_2} (1 - f_i^{\mu_2}) \delta(\varepsilon_i^{\mu_2} - \varepsilon_i^{\mu_1} + \varepsilon_i^{\mu_2} - \varepsilon_i^{\mu_1})
\]

where the first and second terms in the square brackets correspond to direct and exchange Coulomb scattering, respectively. The matrix elements of the screened Coulomb interaction, \( W_{\mu_2 \nu_2 \mu_1 \nu_1} \), are calculated as in Ref. 8. The \( \delta \)-functions describe energy conservation in the Markov limit. Non-Markovian corrections to the carrier scattering are expected to give only small changes on the timescales considered here, see, e.g., Fig. 2 in Ref. 15 for the influence of non-Markovian corrections for the example of the carrier-phonon interaction.

It is known that the carrier-phonon interaction needs to be treated within the polaron picture in QD systems. As the theoretical framework that is applied here for the light-matter interaction does not allow for an easy inclusion of polaron effects, we use a relaxation time approximation, where the interaction does not allow for an easy inclusion of polaron theoretical framework that is applied here for the light-matter interaction, which correspond to direct and exchange Coulomb scattering, respectively. Therefore, we use the rates \( S_i^{\text{in, out}} \) for the carrier-carrier Coulomb interaction as approximations for the rates \( S_i^{\text{in, out}} \) for the carrier-phonon interaction. The former is also influenced by the QD level spacing, with a larger (smaller) level spacing in general leading to lower (higher) scattering rates.

The fast switch-on is enabled by an ultra-fast carrier dynamics, which is seen in the inversion that is shown as the dotted line in Fig. 1.

The steady state value that is reached is mostly controlled by the Q-factor, associated with the photon loss rate. In contrast, it is the combination of carrier kinetics and Purcell factor that determines the slope of the photon number vs. time curve and thus the time needed to reach steady state. The former is also influenced by the QD level spacing, with a larger (smaller) level spacing in general leading to lower (higher) scattering rates.

The fast switch-on is enabled by an ultra-fast carrier dynamics, which is seen in the inversion that is shown as the dotted line in Fig. 1.
The inversion reaches the maximum value after about 3 ps, which points towards extremely efficient carrier capture and relaxation. We also observe that there is little spectral hole burning (shb), as the inversion has a steady state value of \( \approx 0.45 \). This can be explained by the low photon number that is present in this device. Since the average photon number is only about 2 in steady state, the stimulated emission rate is not much larger than other rates in the system, as it would be in conventional laser devices, which would cause significant shb. In contrast, the rates for emission and carrier scattering for such a nanolaser device are of comparable order of magnitude, which means that the degree of shb observed depends on the details of the system. It should be noted that the switch-on behavior cannot be reproduced by using constant scattering times for the Coulomb interaction.

The details of the carrier dynamics are shown in Fig. 2, where the carrier population functions is shown as a function of energy for different times. After the initial Gaussian population profile excited by the pump, we observe an ultra-fast redistribution that immediately starts to populate the QD states. Remarkably, the timescales between the relaxation dynamics in the WL and the dynamics of the QD populations are not decoupled, as often assumed. On one hand, this is due to the efficiency of the initial in-scattering into the QD, as the QD states are initially empty. On the other hand, the relaxation within the WL is significantly slower than in a pure quantum well. The latter effect is caused by the fact that the capture into the QD states is most efficient for the WL states with low quasi-momenta. Therefore, these states are constantly depleted during the early stage of the kinetics, which slows down the relaxation of the WL distribution itself towards a quasiequilibrium distribution. Moreover, a common chemical potential for QD and WL states is reached, as it is expected from experimental data for temperatures over 250 K.

Furthermore, after the initial relaxation that lasts about 3 ps, a heating of the carrier population is observed on a time-scale of 100 ps, as the pumping into the WL injects carriers at a higher energy than the emission energy. Even though the lattice temperature is held at 300 K, the resulting steady-state Fermi distribution has a carrier temperature of \( \approx 1120 \) K. This effect has been observed in QW laser structures, but is more pronounced in QDs, as the energy difference between injected and recombining carriers is larger and has also been discussed for QD devices.

In Fig. 3, we present the non-equilibrium scattering rates as functions of time. After an initial rise, the out-scattering rates both for electrons and holes exhibit a maximum at about 600/ps. The high rate is outbalanced by the low QD occupation (cf. Eq. (4)) during the early stage of the kinetic. After this fast initial rise, the out-scattering rates are reduced significantly and rise again only weakly in conjunction with the above mentioned heating. The in-scattering rates also increase strongly within the first picosecond, leading to the observed fast carrier dynamics. Moreover, there is a sharp increase in the in-scattering rate for the electrons, approximately at the time when the WL distribution for electrons first acquires significant values at the band edge.

In conclusion, we have demonstrated a fast switch-on of QD-based nanolaser devices that is driven by an ultra-fast carrier dynamics. The associated scattering rates exhibit a very strong time dependence, whose inclusion is vital to describe the switch-on behavior. The combination of fast switch-on dynamics and weak relaxation oscillations shows that such laser structures are promising candidates for ultra-compact on-chip signal sources.

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\[ \text{References} \]