Quantum cascade lasers (QCLs) boast a combination of nonlinear dynamic properties that is unique among solid-state lasers. On one hand, an ultrafast (~1 ps) electron relaxation time in excited subbands and an even faster relaxation of the intersubband coherence make QCLs the only solid-state lasers belonging to the group of class-A lasers [1], in which population inversion and polarization adiabatically follow the dynamics of the electromagnetic field. On the other hand, a very large dipole moment of the intersubband transitions in combination with high intracavity intensity gives rise to a large Rabi frequency of the order of 1 THz. Therefore, one can expect that nonlinear effects originating from spectral and spatial hole burning that are only marginal or nonexistent in other types of semiconductor lasers will be ubiquitous in QCLs. Indeed, in recent papers [2,3], we reported observations of multi-longitudinal-mode instabilities including the coherent Risken-Nummedal-Graham-Haken (RNGH) instability in different types of QCLs. In this Letter we present evidence that the nonlinear coupling between the transverse modes of a QCL is equally ubiquitous and gives rise to strong phase coherence effects.

Experimental evidence for beam steering has been recently reported in the far field of high power QCLs and explained by interference of the two lowest order transverse modes [4]. Previously, similar phenomena were observed in diode lasers and also explained as transverse-mode frequency and phase locking; see, e.g., [5–10], and references therein. However, most of the above works dealt with diode lasers and considered only two-mode coupling. The salient features of diode lasers and, in particular, the mechanisms of phase locking are quite different from the relevant properties of QCLs. In diode lasers, the locking is primarily due to density modulation of free carriers, which is proportional to the linewidth enhancement factor \( \alpha \). In QCLs the total electron density is not affected by the laser field and the \( \alpha \) factor is negligible. In this case the mode coupling is due to gain saturation and the concomitant strong spatial hole burning, which favors multimode operation [2,3].

In this Letter, we perform systematic spectral, near-field and far-field transverse-mode measurements of QCLs. We observe phase coherence between multiple transverse modes and explain it by a four-wave mixing interaction. We find that the condition of phase locking is in good agreement with the prediction of the theory.

Buried heterostructure QCLs (\( \lambda \approx 7 \mu m \)) [11] with the same quantum well design but different ridge width (8, 12, 17, and 22 \( \mu m \)) were employed in these measurements. A minimum of two devices for each ridge width was studied. Except for the 8 \( \mu m \) wide QCLs, which were too narrow to operate on more than one transverse mode, all lasers showed clear and consistent evidence of phase locking of multiple transverse modes. In the present Letter, we focus on the results obtained with lasers of two different active region width: 12 and 22 \( \mu m \). High-resolution laser spectra (resolution 0.125 cm\(^{-1}\)) of the 12-\( \mu m \) devices show that there is more than one set of longitudinal modes [Fig. 1(a)], indicating excitation of several transverse modes. A Fourier analysis of this spectrum reveals three closely
spaced peaks [Fig. 1(b)]. It is natural to attribute them to three combs of longitudinal modes belonging to three dominant transverse modes with slightly different effective refractive indices. Similar spectra were observed for lasers with a broad 22-μm active region.

QCL far-field mode profiles were measured using a setup that is based on a rotation stage. Near-field mode profiles were characterized by a midinfrared apertureless near-field scanning optical microscope (mid-ir a-NSOM) [12]. In the experiment, a metal coated atomic force microscope tip oscillates at frequency $f_o$, and scans the laser facet, scattering the near field of the laser mode into the far field. We demodulated the detected scattered light at $2f_o$, in order to discriminate the scattered near field from the background [12]. In order to enhance detection sensitivity, the interference between the near field scattered by the tip apex and the direct laser output is detected, and the NSOM image is a map of the amplitude of the near field of the transverse mode [12,13].

The measured far-field and near-field mode profiles for two sets of devices with widely different ridge widths (12 and 22 μm) show clear evidence of coherent coupling of multiple transverse modes [Figs. 2(b), 3, 4(b), and 5]. To extract information of the constituent transverse modes from the measured mode profiles, we first perform simulations to solve the transverse eigenmodes of the QCL waveguide. Then a near-field mode profile is expressed as a combination of several transverse eigenmodes including their relative amplitudes and phases. A Fresnel diffraction integral is then used to calculate the corresponding far-field mode profile. By fitting it to the experimental far-field profiles, we find that there are three constituent transverse modes with specific phase and amplitude relations [Figs. 2(c) and 4(c)]. The far-field fits are quite sensitive to small changes of the phases and amplitudes of the constituent transverse modes. In order to quantify this, we use error bars in Figs. 2(c) and 4(c) to indicate the range of individual parameter changes that could lead to a root mean square deviation of 0.05 between the experimental data and the fits. We note that under small driving current conditions, we cannot reliably determine the amplitudes and phases of the two weak transverse modes. The uncertainty of fitting parameters for small current conditions therefore is large [Figs. 2(c) and 4(c)]. However, the small current regime is not of interest to us; our study focuses on laser operation at sufficiently large currents when the intensity of the three modes is high enough to ensure strong nonlinear coupling.

As an independent check of the reliability of the far-field fits, we plot the experimental a-NSOM near-field mode profiles together with the near-field profiles obtained using the correct parameters determined from far-field fits. We find a reasonably good agreement between the two [Figs. 3(b) and 5(b)]: the coefficient of determination $R^2$ for the near-field fits was found to be about 0.95 for the 12-μm wide device and about 0.90 for the 22-μm wide device. The correlated far-field and near-field measure-
ment results form solid experimental evidence of phase locking.

We propose the following simple model of phase coherence between transverse modes. In the presence of a Fabry-Perot laser cavity, each transverse mode of a waveguide consists of a comb of longitudinal cavity modes. At low laser intensities, longitudinal modes belonging to different transverse modes are uncorrelated. At high intensities, however, these longitudinal modes become coupled due to four-wave mixing mediated by the saturation nonlinearity of the laser transition. We show that the latter process leads to the frequency and phase locking between longitudinal modes belonging to different combs. Our theory predicts definite values for the relative amplitudes, frequencies, and phases of locked modes that turn out to be in excellent agreement with observations.

Note that there is no contradiction with the known result stating that four-wave mixing of longitudinal modes in the two-level active region alone cannot lead to a stable phase-locking [1]. This result is not applicable when interacting longitudinal modes belong to different transverse modes. Feasibility of a stable phase locking for two transverse modes was shown, e.g., in Ref. [4] as well as in Ref. [10] and Refs. [1–5] therein. However, the modal frequencies have to be locked to the same values. This is not the case with our lasers where multiple transverse modes at different frequencies are clearly distinguishable in the spectra.

We start from standard Maxwell-Bloch equations [1,2] in which the electric field in a laser cavity is expanded into a series over a set of the cavity modes \( E_j(r) = c_1 \sigma_1 E_j(r) + c_c \), where \( \sigma_0 \) is a central frequency of the laser transition. Since relaxation times \( T_1 \sim 1 \) ps and \( T_2 \sim 0.1 \) ps of the population inversion \( N \) and the coherence \( \rho_{21} \) (i.e., the off-diagonal element of the density matrix) of the laser transition are much shorter than the photon lifetime \( \sim 10 \) ps, we can adiabatically eliminate \( N \) and \( \rho_{21} \) and arrive at the equations for the field amplitudes in the \( \chi^3 \) approximation:

\[
d a_j/dt + (\alpha_j + i(\omega_j - \omega_0))a_j = \frac{g_j}{2} \left\{ \sum_k G_{jk} dV \int_{AR} eE_kE_m dV - \frac{1}{I_s} \sum_{l,m} G_{jklm} a_k^* a_m + a_l a_l^* \right\}. \tag{1}
\]

where the overlap integrals \( G_{jklm} = \int_{AR} eE_jE_kE_m dV \) over the active region volume are calculated numerically, \( g_j = 4\pi\omega_0^2 N_p T_2/(\hbar \mu_j^2) \) is the small-signal material gain, \( \mu_j \) the effective refractive index of the \( j \)th mode, \( I_s = 2\hbar^2/(d^2 T_1 T_2) \) the saturation intensity, \( \omega_j \) and \( \alpha_j \) are the frequency and the modal loss of the \( j \)th “cold” mode, \( d \) the dipole moment of the laser transition, and \( N_p \) the population inversion supported by pumping in the absence of laser generation. Here cold denotes modal parameters calculated assuming no interaction between modes.

The second term on the right-hand side of Eq. (1) describes the phase-sensitive nonlinear coupling between different modes which can lead to their frequency and phase locking. Since our lasers have three transverse modes involved in generation, we divide the whole spectrum of these overlapping modes into triplets of longitudinal modes, with each mode in the triplet belonging to a different transverse mode 1, 2, or 3, and consider the four-wave mixing process \( j + j \leftrightarrow k + l \) within each triplet. We express the complex amplitude of each longitudinal mode via its real amplitude and phase as \( a_j(t) = f_j(t) \times \exp(i(\varphi_j(t) - \Delta_j t)) \), where \( \Delta_j = \omega_j - \omega_0 \) is the detuning of the frequency of a \( j \)th “hot” mode from the transition frequency. Here hot denotes actual (unknown) frequencies, phases, and amplitudes of laser modes with the nonlinear interaction taken into account. Then Eq. (1) yields six equations with real coefficients for the amplitudes and phases of these modes. For 12-\( \mu \)m wide lasers, three transverse modes 1, 2, and 3 are TM\(_{01}\), TM\(_{01}\), and TM\(_{02}\) modes, respectively. We find that the process that can lead to phase locking is of the type TM\(_{01}\) + TM\(_{01}\) ↔ TM\(_{00}\) + TM\(_{02}\), which corresponds to the phase combination 2\( \varphi_2 \) − \( \varphi_1 - \varphi_3 \):

\[
d (2\varphi_2 - \varphi_1 - \varphi_3)/dt + \delta_1 + \delta_3 - 2\delta_2 = -(g/12)(G_{1222}(2f_1^2 + f_3^2)\sin 2\Phi_{12} + 2G_{2333}(2f_3^2 + f_3^2)\sin 2\Phi_{32} + G_{1133}(f_1^2 - f_3^2) \times \sin 2\Phi_{31} + 2G_{1233}(f_1^2 - f_3^2)\sin 2\Phi_{12} + G_{1223}(f_3^2 - f_3^2)\sin 2\Phi_{31} + 2G_{1233}(f_3^2 - f_3^2)\sin 2\Phi_{12} + G_{1133}f_3^2\sin 2\Phi_{31}). \tag{2}
\]

In Eq. (2) \( \delta_j = \omega_j - \omega_{ej} \), \( \Phi_{jk} = \varphi_j - \varphi_k - (\omega_j - \omega_k)t \), and \( \Phi_{2333} = 2\varphi_2 - \varphi_1 - \varphi_3 - (2\omega_2 - \omega_1 - \omega_3)t \).

Phase locking \( d (2\varphi_2 - \varphi_1 - \varphi_3)/dt = 0 \) occurs when three conditions are satisfied: \( f_1 = f_3 \); the frequency locking of the hot modes \( \omega_1 + \omega_3 = 2\omega_2 \); the mode amplitudes are sufficiently large so that the nonlinear term on the right-hand side of Eq. (2) cancels the term \( \delta_1 + \delta_3 - 2\delta_2 = 2\omega_2 - \omega_1 - \omega_3 \) originating from the waveguide dispersion of the cold modes. These conditions agree with the observed behavior of the amplitudes, phases, and frequencies of the transverse modes. Indeed, as one can see from Figs. 2(b) and 2(c), the lasing starts on the TM\(_{01}\) mode which has the lowest threshold [12]. With increasing current, the TM\(_{00}\) and TM\(_{02}\) modes become involved in generation and become nonlinearly coupled with the TM\(_{01}\) mode by the four-wave mixing process. Above the locking-threshold current of the order of 0.8 A, the amplitudes of the TM\(_{00}\) and TM\(_{02}\) modes become very close to each other.
and the phase difference $2\varphi_2 - \varphi_1 - \varphi_3$ does not change with current [Fig. 2(c) lower panel]. Furthermore, at low currents, the three peaks in the lower panel of Fig. 1(b), corresponding to three combs of the longitudinal modes, are nonequidistant because of the cold cavity dispersion. Above the locking-threshold current, the triplet of longitudinal modes belonging to three different combs are locked into the equidistant relationship $\omega_2 - \omega_1 = \omega_3 - \omega_2$. This results in three equidistant peaks in the upper panel of Fig. 1(b).

One can perform a similar analysis for 22-µm wide lasers that start lasing on the TM$_{00}$ mode, which has the lowest threshold [12]; see Figs. 4(b) and 4(c). With increasing current, the modes TM$_{00}$ and TM$_{02}$ become involved in generation. Therefore, three dominant coupled transverse modes 1, 2, and 3 for these devices are TM$_{00}$, TM$_{02}$, and TM$_{04}$ modes, respectively. We find that the process that can lead to phase locking is of the type $2\varphi_3 - \varphi_1 - \varphi_2 = 2\pi n$, with $n = 0, \pm 1, \ldots$, when $f_2 = f_3$, $\omega_1 + \omega_3 = 2\omega_2$, and the modal amplitudes are high enough to beat the dispersive term $\delta_1 + \delta_2 - 2\delta_3 = 2\omega_3 - \omega_1 - \omega_2$. The above conditions are again in excellent agreement with observations.

In conclusion, using near-field and far-field measurements we observed phase coherence between several transverse modes in QCLs over a broad range of injection currents. Our data can be explained by phase-sensitive nonlinear coupling of transverse modes due to resonant $\chi^{(3)}$ nonlinearity of the active medium. Evidence for stable transverse mode locking was observed for all 12, 17, and 22-µm wide devices processed from the same wafer. We believe that transverse mode locking is a universal phenomenon for high power QCLs that support multiple low-loss transverse modes. The resulting frequency- and phase-locked multimode spectrum can be further utilized for creating various far-field radiation patterns, control of the beam quality and beam steering, and facilitating ultrashort pulse generation.