# Acoustic Phonon Emission from a Weakly Coupled Superlattice under Vertical Electron Transport: Observation of Phonon Resonance

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We report measurements of acoustic phonon emission from a weakly coupled AlAs/GaAs superlattice (SL) under vertical electron transport. The phonons were detected using superconducting bolometers. A peak (resonance) was observed in emission parallel to the SL growth axis when the electrical energy drop per SL period matched the energy of the first SL mini-Brillouin zone-center phonon mode. This peak was mirrored by an increase of the differential conductance of the SL. These results are evidence for stimulated emission of terahertz phonons as previously predicted theoretically and suggest that such a SL may form the basis of a SASER (sound amplification by stimulated emission of radiation) device.

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Recently it was predicted that for a weakly coupled semiconductor superlattice (SL) in the hopping vertical electrical transport regime, the conditions necessary for terahertz phonon amplification can be achieved [1]. Possible experimental evidence for such amplification has been observed in phonon-assisted transport measurements on GaAs/AlAs SLs irradiated by an external source of nonequilibrium phonons [2,3]. Here we report the observation of resonancelike emission of terahertz acoustic phonons, which occurs when the Stark splitting of the adjacent quantum-well (QW) levels matches the energy of miniband-center acoustic phonons. This is exactly the condition when efficient feedback is realized for phonons with high amplification, which is qualitatively similar to distributed feedback (DFB) lasers. This observation suggests the possibility to develop an electrically pumped high-intensity source of terahertz coherent acoustic phonons based on a SL structure. Such a device, which could be called a SASER (sound amplification by stimulated emission of radiation), would have potential applications in phonon optics, phonon spectroscopy, and acoustical imaging of nanostructures [4].

In SL structures, the artificial periodicity of the lattice potential along the growth (or $z$) direction gives rise to folding of the conduction band into a series of minibands separated by minigaps [5]. The miniband width, $\Delta_m$, depends on the probability of tunneling between adjacent QWs and is narrow if they are only weakly coupled. Subject to an applied electric field, $F$, such that $eFd_{\mathrm{SL}} > \Delta_m$, where $d_{\mathrm{SL}}$ is the SL period, the electron states become strongly localized within individual QWs. Vertical electron transport then takes place by tunneling and electrons "hop" between states in adjacent QWs. Conservation of energy and momentum requires that such transport takes place either by an elastic process, involving defect scattering with the subsequent emission of phonons, or by inelastic phonon-assisted tunneling, see Fig. 1. In Ref. [1] it was argued that the initial and final states for transitions involving emission of phonons close to the $z$ direction can be population inverted (see bottom part of Fig. 1). Under such conditions, phonon amplification by stimulated emission is possible. Experimental measurements of acoustic phonon-assisted tunneling in a GaAs/AlAs superlattice [2,3] showed that, for a given value of $F$, an incident phonon pulse gave rise to an increase, $+\Delta I$, in the electron current due to tunneling. The value of $F$ at which $\Delta I$ was largest

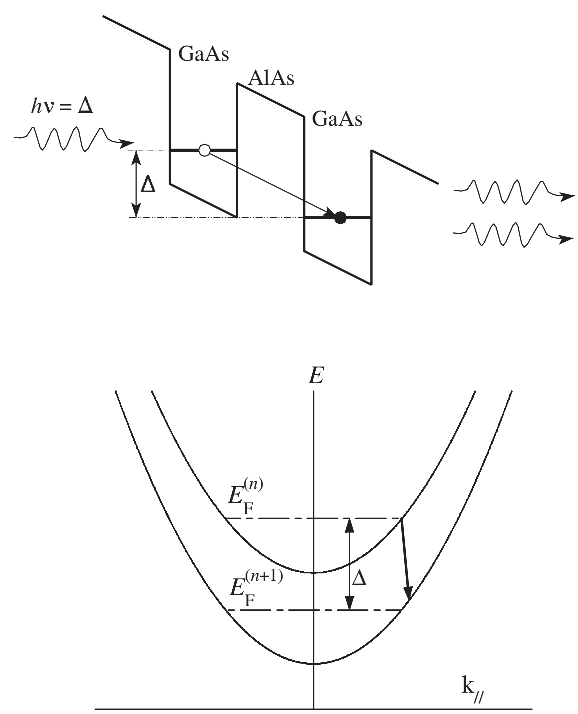


FIG. 1. Illustration of phonon-assisted (via stimulated emission) hopping conduction between a pair of neighboring GaAs wells in a SL. In the bottom portion of the figure the 2D subbands corresponding to 2D electrons confined in the adjacent QWs are shown. The indirect phonon-assisted electron transitions characterized by the population inversion are indicated.

 

was found to depend on the phonon energy. It was noted that, at small angles of incidence of the phonon pulse with respect to the SL growth direction, the magnitude of $\Delta I$ was somewhat larger than expected from theoretical estimates using the two-well model [6], and that this might be due to phonon amplification taking place. If an acoustical cavity is created in the system, phonon losses are minimized and the phonon feedback can give rise to generation of monochromatic nonequilibrium phonons [7].

The simplest way to achieve the feedback is to use modification of the phonon properties in the SL. As well as giving rise to folding of the conduction band, the periodic nature of the SL also leads to folding of the acoustic phonon dispersion. This is due to the different materials which make up the SL having different acoustic impedances. Narrow gaps open up at the phonon energies where the folded dispersion meets the minizone center and minizone boundary. The phonon modes near these gaps are confined in the SL. This means that the SL can act as its own acoustic cavity, analogous to a DFB laser cavity. All that is needed is to ensure that the gain frequency of the SL coincides with the cavity frequency. This can easily be done by applying the appropriate value $F$ to the SL.

The experimental device, Fig. 2, was based on a 50-period GaAs/AlAs SL grown by molecular beam epitaxy on a 0.4 mm-thick semi-insulating GaAs substrate. Each period consisted of a 5.9 nm-thick GaAs well and a 3.9 nm-thick AlAs barrier, uniformly $n$ doped with Si to a density of $2 \times 10^{22}$ m$^{-3}$. The SL was separated from the $n^+$ ($2 \times 10^{24}$ m$^{-3}$) contact regions by 20 nm-thick undoped GaAs spacer layers. A 50 $\mu$m-diameter device mesa was formed by etching and contacts with the emitter and collector layers made using GeAuNiAu, alloyed at 360 °C. Using the Bastard model [8], we estimate the miniband width in this SL to be 0.7 meV. The device was characterized at temperature $T = 1.5$ K by measuring its dc current-voltage ($I$-$V$) characteristics, also shown in Fig. 2. The device turns on at the threshold bias, $V_T \simeq 75$ mV, which is the voltage required to align the Fermi energy of the emitter and the nearest well. After that, the current increases monotonically for biases up to about 250 mV. Fluctuations of the current with applied bias above 250 mV indicate the formation and growth of electric field domains within the sample [5]. For measurements made at bias voltages between $V_T$ and 250 mV it is assumed the variation of electric field along the structure is uniform. In this case, the energy drop per period of the SL is given by $\Delta = \gamma e(V - V_T)/50$, where $\gamma$ is the fraction of the applied voltage that is dropped across the SL and is determined empirically from the phonon-assisted tunneling measurements, as described in detail in Ref. [2]. For this sample under negative bias to the top contact, as used here, $\gamma \simeq 1$.

The back surface of the substrate was polished and a pair of $40 \times 40$ $\mu$m$^2$ active-area superconducting aluminum bolometers were fabricated. One was positioned directly opposite the device, the second subtended an angle of $\theta = 30°$ to the $z$ direction, see Fig. 2. All measurements were made with the sample in a liquid helium cryostat at temperature $T \approx 2$ K, on the superconducting transition of the bolometers. Voltage pulses of duration 1 $\mu$s and amplitude in the range 0–0.5 V were applied to the device. The emitted acoustic phonons propagated ballistically across the substrate and were incident on the bolometer. The transient signal from the bolometer was amplified and detected using a fast digitizer and signal averager.

Figure 3(a) shows the folded longitudinal acoustic (LA) phonon dispersion for the SL described above. For phonons propagating in the $z$ direction, the gaps are centered on phonon frequencies, $\nu_n = \frac{n}{2}(\frac{d_{\mathrm{GaAs}}}{c_{\mathrm{GaAs}}} + \frac{d_{\mathrm{AlAs}}}{c_{\mathrm{AlAs}}})^{-1}$, where $d$ and $c$ are, respectively, the thickness of and sound velocity in the appropriate SL layer. In our SL, the first and second minizone edge gaps are at $\nu_1 = 256$ GHz (boundary) and $\nu_2 = 512$ GHz (center), respectively. The acoustic reflection coefficient of the SL, calculated following the method described in [9], is shown in Fig. 3(b).

The dependence on $\Delta$ of the detected phonon intensity is shown in Fig. 4. This result was obtained using the bolometer directly opposite the device. The intensity has been normalized to the power dissipated in the device, which was determined from measurements of the current-voltage characteristics. The bolometer is an energy detector and so,

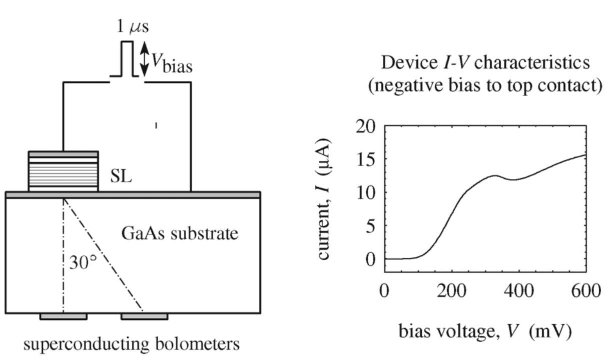


FIG. 2. Experimental sample structure and $I(V)$ characteristics of the SL device.

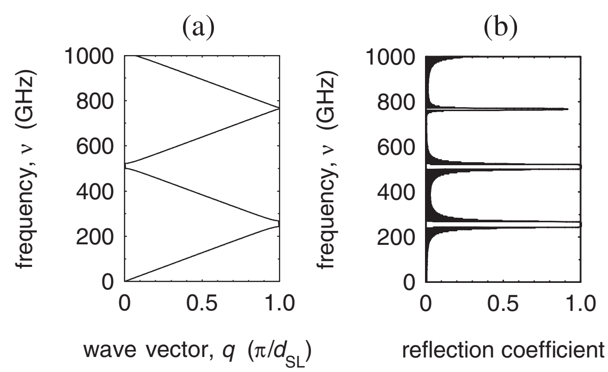


FIG. 3. (a) Folded phonon dispersion calculated for the SL used in the measurements; (b) acoustic reflection coefficient (as a function of phonon frequency) for the SL.

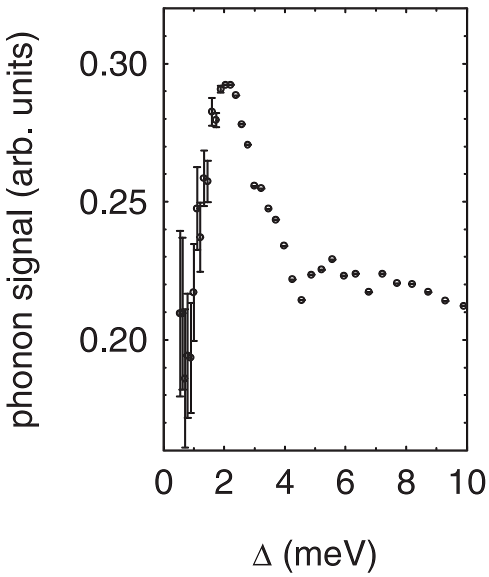


FIG. 4. Phonon signal, normalized by the total power dissipated in the device, as a function of $\Delta$.

if the phonon intensity was simply proportional to the power dissipated, we should expect a featureless horizontal line. Instead we see a clear broad peak centered on $\Delta \simeq 2.1$ meV. No such peak, just a featureless horizontal line, was observed using the bolometer at $\theta = 30°$. The fact that we observe a peak in the normalized signal suggests that there must be a change occurring in the spatial or spectral distribution of the phonon emission at $\Delta \simeq 2.1$ meV resulting in an increased phonon flux reaching the detector. Possibilities are increased emission in the $z$ direction at the expense of other directions and/or increased emission of low energy (sub-THz) acoustic phonons which suffer less scattering on their way to the detector.

First we should dismiss the most obvious explanation of the peak in the normalized phonon signal: that it is simply due to an increase in the device dissipation at $\Delta \simeq 2.1$ meV. The absence of any strong features in the power dissipation near $\Delta = 2.1$ meV would suggest this not the case. However, the normalization should negate any such effect and, furthermore, no peak was observed using the bolometer at 30°. Additionally, the angle and frequency distribution of the phonons emitted due to carrier relaxation in the device (including the contacts) is determined by energy and momentum conservation and should not show a strong bias dependence.

A hint as to a possible explanation is provided by the value of $\Delta$ at which the peak is observed: phonon-assisted transitions between adjacent QWs separated in energy by 2.1 meV requires phonons of frequency 510 GHz, almost exactly the same as the lowest zone-center SL mode for LA phonons. To put it another way, the peak in the phonon emission normal to the SL layers is observed when the Stark splitting between adjacent QWs is resonant with a confined phonon mode in the SL. This suggests the increased emission could be as a result of phonon amplification occurring in the device coupled with the SL behaving as an acoustic "cavity" providing feedback. According to [1], phonon amplification is possible for phonons traveling at angles $\theta < \theta_{\max} = \sin^{-1}\frac{c}{v_F}$, where $v_F$ is the Fermi velocity of the electrons in the QWs. For our device, $v_F \approx 6 \times 10^4$ ms$^{-1}$; therefore, $\theta_{\max} \approx 5°$. This might explain why no peak in the emission with changing $\Delta$ was observed using the bolometer at $\theta = 30°$. It is worth pointing out that the broad shape of the observed peak is not in contradiction to the resonant character of the phonon feedback. Even in an ideal SL, generation is predicted in a relatively wide range of Stark splittings where the phonon gain is sufficiently large. In real structures further broadening is possible due to variation of the layer parameters.

Another possibility that must be considered is that the SL has a strongly frequency dependent emissivity for phonons which, when convolved with the spectral distribution of spontaneously emitted phonons of energy $\Delta$, leads to the observed peak. We believe this is not the case because within the phonon stop bands, the acoustic reflection coefficient of the SL is large. By analogy with multilayer optical mirrors, we see that this should give rise to a reduction in the emissivity and a corresponding reduction in the emission, not the increase observed.

If the peak in emission in the $z$ direction is due to phonon oscillation in the structure, then there should be an increase of phonon occupation at the energy required for phonon-assisted hopping. This should influence the SL current-voltage characteristic. In particular, a peak in the differential conductance, $dI/dV$, of the device at the appropriate value of $\Delta$ is expected. We measured the differential conductance of the device using a small ac (1 mV; 37.5 Hz) modulation of the bias and phase sensitive detection of the current. The results are shown in Fig. 5. A clear peak in $dI/dV$ is observed at $\Delta = 2.1$ meV which is exactly where the peak in the phonon emission was observed. It is inter-

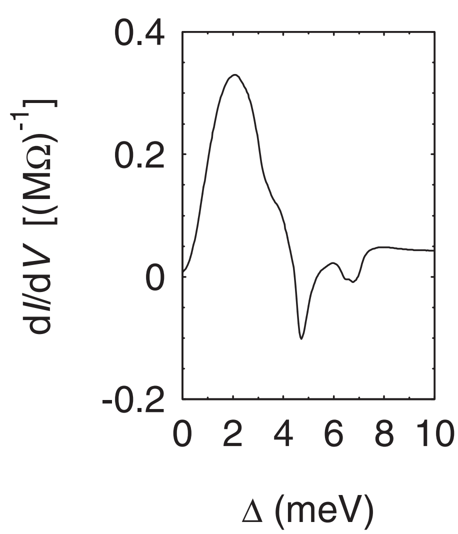


FIG. 5. Differential conductance of the SL device as a function of $\Delta$.

esting to note the close similarity between the normalized bolometer signal shown in Fig. 4 and $dI/dV$, even with regard to the minor features occurring at higher bias than the main peak. It is, perhaps, not surprising that the conductance mirrors the phonon emission in this way because electron transport takes place either by direct phonon-assisted hopping or by elastic scattering followed by energy relaxation, and both involve phonon emission. We can discount the possibility that this similarity is due to an experimental limitation related to charging of the device because we deliberately used a long (1 $\mu$s) pulse and made the bolometer measurement only after the device had time to reach the steady state. From Fig. 5 we estimate that the generated phonon power is of the order milliwatts per cm$^2$, which is considerably less than the values predicted in [7]. This means there is scope for improving the efficiency of generation by more accurate design of the SL structure.

At first it might appear that these results are in contradiction to the theory developed in [1] which predicts no emission for miniband-center phonons. However, the origin of this prediction is in the particular model of hopping transport used rather than creation of the population inversion. The theoretical model assumed hybridization of the electron orbitals belonging to the adjacent QWs, which corresponds to the so-called Wannier-Stark hopping conduction regime [10]. This leads to the approximate expression for the wave function corresponding to electrons confined mainly in the $i$th QW: $\Psi_i = \chi_i + (t/eFd) \times (\chi_{i-1} - \chi_{i+1})$, where $\chi_n$ are orbitals corresponding to $n$th QW and $t$ is the tunneling matrix element (in this approximation is $\Delta_m = 4t$). Consequently, the probability of phonon-assisted interwell transitions is proportional to the form factor $J_{\text{inter}} \sim \sin^2(q_z/2d) J_{\text{intra}}$, where $q_z$ is the $z$ component of the phonon wave vector and $J_{\text{intra}}$ the form factor for intrawell phonon-assisted transitions. It is that property which suppresses emission of miniband-center phonons. This picture is destroyed if the electron scattering time $\tau < \hbar/t$. This corresponds to the case of sequential tunneling [10], where hybridization is destroyed by strong scattering. Here, the interwell form factor does not contain the factor $\sin^2(q_z/2d)$, and emission of miniband-center phonons is allowed. This regime of transport, as in the case of Wannier-Stark hopping, is characterized by an inverted electron distribution. For the structure under investigation, the effective-mass estimate gives $t \sim 0.25$ meV and, based on measurements of single-QW structures grown in a similar way and with similar doping, we estimate the electron mobility to be about $5 \times 10^3$ cm$^2$ V$^{-1}$ s$^{-1}$. According to these estimates, $\hbar/\tau \sim 3$ meV; therefore, the structure under investigation is characterized by sequential tunneling and so there is no restriction on the emission of miniband-center phonons.

One remaining question that needs to be addressed is why resonance with the first SL minizone-center phonon mode dominates the emission and not, for example, the first or second minizone boundary modes at 256 GHz ($\Delta =$ 1.05 meV) and 768 GHz ($\Delta = 3.15$ meV), respectively? There is a clear feature in the normalized bolometer signal at $\Delta \approx 3$ meV which could be due to resonance with the second minizone boundary mode. From Fig. 3(b) one can see that the confinement of such phonons is relatively weak. Roughly, the quality factor of confined phonons can be estimated as $Q \sim 1/(1 - r)$, where $r$ is the phonon reflection coefficient. For the second zone boundary phonon we have $1/(1 - r) \approx 12$, while for the first minizone center $1/(1 - r) \approx 10^5$, which probably explains the difference in phonon emission at $\Delta = 3.15$ meV and $\Delta =$ 2.1 meV. On the other hand, no convincing feature is visible at $\Delta \approx 1$ meV. This could be due to weak electron confinement in individual QWs at such small bias: the estimated miniband width in our sample is about 1 meV, which is close to the corresponding bias. This prevents localization of electrons in individual QWs and formation of population inversion in the system.

In summary, we have observed resonantlike emission of acoustic phonons from a weakly coupled SL in the vertical hopping transport regime. We attribute this emission to acoustic phonon amplification supplemented by the phonon feedback due to stop bands in the SL phonon spectrum which results in efficient phonon generation. Phonon generation is accompanied by a peak in the SL differential conductance which allows us to estimate the phonon power to be in the range mW cm$^{-2}$.

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