Manipulation of graphene’s dynamic ripples by local harmonic out-of-plane excitation

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Abstract

With use of carefully designed molecular dynamics simulations, we demonstrate tuning of dynamic ripples in free standing, thermally fluctuating graphene by applying a local out-of-plane sinusoidal excitation. The local dynamic morphology can be controlled via varying external modulation and the boundary conditions. We fully account for the discrete atomistic structure of graphene, as well as natural energy dissipation due in part to its remarkably high thermal conductivity. In addition to stable dynamic rippling patterns, we observed an unexpected flattening of graphene well below the thermal limit. Our results provide insight into the dynamic response of atomically thin layers to an external time-varying excitation in the presence of realistic thermal fluctuations and energy loss. This suggests intriguing possibilities for modulating the electrical and optical properties of atomically thin membranes via local dynamic morphology control.

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(Some figures may appear in colour only in the online journal)

1. Introduction

Atomically thin layers present a unique case in solid-state physics, where two-dimensional confinement results in remarkable electronic properties of graphene [1–4], the first truly two-dimensional (2D) crystal [1, 5]. The effect of local landscape of the layer on its properties is important, because free-standing graphene is naturally rippled [6]. These ripples have been attributed to a dynamic effect of finite temperature [7], ultimately resulting in the destruction of the long-range order in 2D crystals, in qualitative agreement with the long-standing Mermin–Wagner theorem [8] and the statistical considerations of the thermally fluctuating membranes [9, 10]. This mechanism is closely related to the fundamental instability of the Bose–Einstein statistics in strict two dimensions [11] and can be viewed as graphene’s natural way of existing in three-dimensional space. Static rippling has been suggested as an important additional mechanism affecting the local morphology, arising from disrupted bonding coordination at the edges, as well as from lattice defects dispersed throughout the layers [12, 13]. Static wrinkles in substrate-bound graphene have also been reported, along with their effect on graphene’s electronic properties [14].

Control of the texture in suspended graphene sheets by local strain [15], due to rolling [16] and dynamic mechanical excitation of graphene [17], was previously reported. These findings are especially important, because recent experimental work revealed local quantum confinement in graphene drumheads [18], which results from adjustable strain-induced pseudomagnetic fields previously predicted theoretically [19]. In light of these reports, our ability to control the local morphology of suspended atomically thin layers may result in a new class of mechanically controlled quantum dots.

Manipulation of the local dynamic shape of suspended atomic layers can be achieved via dynamic time-varying
excitation, similar to the case of vibrating drumheads. Here, we report the results of molecular dynamics (MD) simulations of an out-of-plane sinusoidal excitation locally applied to a thermally fluctuating graphene membrane. We demonstrate that the coexistence of stochastic (thermal) ripples and the externally excited out-of-plane vibrations in an atomically thin membrane subject to energy dissipation results in a number of interesting phenomena, including significant overall flattening and emergence of stable dynamic ripple patterns. We show high controllability of the wavelength and amplitude of the dynamic membrane ripples, depending on the excitation amplitude and frequency, beyond the previous results for electromechanical graphene resonators [17].

2. Simulated system

The suspended graphene membranes (dimensions 13.6 nm × 15.8 nm, 8192 atoms) were positioned in the XY-plane and set up as shown in figure 1. In-plane periodic boundaries were imposed to minimize the effect of wave reflections. The equations of motion were integrated in three dimensions with a time step of 1 fs. The interatomic interactions were defined by the Tersoff–Brenner bond-order potential, which is well established in describing a variety of static and dynamic lattice properties of graphene and carbon nanotubes [20, 12, 21–24]. The particular energy form used in this work was described in [24].

Because the thermal fluctuations may be influenced by the simulated heat bath into which the system is immersed, the temperature control in the MD simulations was applied in an extremely careful way. The widely used Nosé–Hoover thermostat [25] was applied to a 1 nm wide region at the perimeter of the membrane, representing the heat sink effect of the sample mount, as denoted by ‘thermostatted atoms’ in figure 1, to ensure that no preset amount of effective viscoelasticity was imposed upon the entire membrane. All averages of interest were calculated for the non-thermostatted ‘free atoms’ (figure 1). These atoms were allowed to move freely, governed only by the interatomic interactions. The velocities at the start of all simulations were generated according to the Maxwell–Boltzmann distribution at \( T_0 = 300 \) K and the temperature maintained at the perimeter with and without external excitation was also 300 K, unless stated otherwise. The effective dissipation was moderate with a relaxation time of 1 ps [26], allowing the thermostat-free region of the membrane to heat up well above \( T_0 \) with sufficiently high external excitation. All simulations were run for 0.2 ns and the averages of interest were calculated for the second half of the simulated time period, determined from the system energy time evolutions as sufficient to reach steady state.

3. Results and discussion

The main quantity of interest in our simulations was the mean square displacement (MSD) of the atoms in the out-of-plane direction (the membrane initially positioned in XY-plane at \( Z = 0 \)), defined as

\[
\langle h_{0t}^2 \rangle = \left( \frac{1}{N} \sum \bar{z}_i^2 \right),
\]

where \( \bar{z}_i \) is the out-of-plane (Z-) coordinate of the \( i \)th atom in the system, \( N \) is the total number of atoms, and \( \langle \cdot \rangle \) represents the time average. Without external excitation, the amount of stochastic out-of-plane rippling (\( \langle h_{0t}^2 \rangle \)) is directly proportional to the equilibrium temperature, as shown in the inset of figure 1, consistent with previous results [9, 7]. At 300 K, we calculated \( h_0 = \sqrt{\langle h_{0t}^2 \rangle} = 0.27 \) Å, which corresponds to an average height of \( \sqrt{3}h_0 = 0.38 \) Å, about an order of magnitude higher than that obtained from the Green’s function based analysis of vibrations at the atomic site [27] and \( \sim 1.8 \) times lower than for a sample of similar size obtained from Monte Carlo simulations [7]. Our thermal ripple height is also significantly lower than the \( \sim 1 \) nm obtained in a continuum model of externally excited graphene membranes [28]. One must keep in mind that the average ripple height scales with the membrane size, as well as any external input and temperature [9], and thus a direct comparison is only reasonable for heights properly normalized with respect to these factors. Within the selected temperature range, there is no indication of significant anharmonic effects, in agreement with [7]. The external excitation was a sinusoidal force \( F(t) = F_0 \cos(\omega t) \) applied to a single atom at the intersection of the diagonals of the membrane, as shown in figure 1. Such input caused the induced vibrations to spread throughout the entire membrane [29] (also see section 1 of the supplementary information available at stacks.iop.org/Nano/24/055701/mmedia).

Shown in figure 2(a) are the results for the root-mean-square of the out-of-plane displacements due to an external force normalized to the corresponding value without external excitation \( \varepsilon = \sqrt{\langle h_{0t}^2 \rangle}/(h_0^2) \). The forcing frequency \( f = \omega/2\pi \) was swept over the indicated range at four different values of \( F_0 \), as shown. For \( F_0 = 8 \) nN and \( f < 5 \) THz, \( \varepsilon \) is considerably
higher than unity and dynamic rippling patterns emerge, as shown in figure 2(b) for \( f = 1 \) THz, while for higher excitation frequencies \( \varepsilon < 1 \), and an intriguing effect of flattening takes place, as also shown in figure 2(b) for \( f = 10 \) THz. The two particular local minima in \( \varepsilon \) are at \( f \approx 11 \) THz and \( f \approx 22 \) THz. Toward the frequency of 40 THz, the membrane appears to lose all its sensitivity to external excitation and \( \varepsilon \) becomes close to unity, while the ripple distribution of the sheet becomes nearly identical to the case without external excitation (figure 2(b)). For \( F_0 = 4, 2, \) and 1 nN, the data in figure 2(a) show nearly identical behavior, although the positions of the minima exhibit shifting to the left. This shifting is likely due to the decreasing anharmonicity in the response with decreasing \( F_0 \). The overall effect on \( \varepsilon \) also decreases with decreasing \( F_0 \), as one may expect. We shall come back to the discussion of the observed dynamic patterns at low frequencies and their significance later in the text.

In order to qualitatively assess the validity of the results described above, we estimate the value of \( \varepsilon \) analytically. The main purpose here is to show that the observed flattening is an important dissipative effect, which can be physically explained by use of thermodynamics. Let us consider the energies associated with the atomic out-of-plane motions in a two-dimensional thermally fluctuating membrane. In the absence of external excitation, the average kinetic energy of the atoms \( E_0 \) (subject to dissipation both in nature and in the simulations) is determined by the effective average temperature of the sheet, calculated as the ensemble average of the per-atom kinetic energy. At the same time, the out-of-plane MSD defined by equation (1) in the absence of external excitation, \( \langle h_0^2 \rangle \), is also proportional to the temperature \([10, 9]\) due to equality between average potential and kinetic energies in the oscillatory steady state:

\[
T_0 \propto E_0 = A \langle h_0^2 \rangle, \tag{2}
\]

where \( A \) is a constant, which depends on the out-of-plane bending rigidity and size of the membrane, arising from the statistical analysis of the thermal ripples [9]. When an external excitation is applied, we assume a harmonic response, i.e. any out-of-plane stochastic (thermal) ripples with MSD of \( \langle h^2 \rangle \) coexist with the deterministic externally induced ripples with MSD of \( \langle h_{ext}^2 \rangle \) (the brackets here represent a time average), and \( \langle h, h_{ext} \rangle = 0 \). As a result of external excitation, as well as any energy loss, the new overall temperature is \( T \); the overall energy is a sum of the thermal energy \( A \langle h^2 \rangle \) and an additional component \( E_{ext} \) carried by the externally induced ripples:

\[
T \propto E = A \langle h^2 \rangle + E_{ext}. \tag{3}
\]

It is paramount to keep in mind that equation (3) is merely a sum of average kinetic energies, where \( A \langle h^2 \rangle \) is the component incorporating all stochastic contributions expressed for convenience as a function of \( \langle h^2 \rangle \) and \( E_{ext} \) is the externally excited kinetic energy contribution. Their distinction is only possible within the harmonic approximation. As we show later, the result we seek does not explicitly depend on the constant \( A \), provided it is identical in equations (2) and (3) (guaranteed by the assumption of harmonic response).

Introducing \( \lambda \) as the general factor of temperature increase as a result of external excitation, as well as any energy loss (\( \lambda \equiv T/T_0 \)), and combining equations (2) and (3), we obtain

\[
\langle h^2 \rangle = \lambda \langle h_0^2 \rangle - \frac{E_{ext}}{A}. \tag{4}
\]

Once again, we note that \( \langle h^2 \rangle \) in equations (3) and (4) is only the stochastic (thermal) portion of the overall out-of-plane vibrations, which cannot be obtained from the simulations. The total amount of simulated rippling defined by equation (1) includes the externally induced contribution \( \langle h_{ext}^2 \rangle \):

\[
\langle h_{ext}^2 \rangle = \langle h^2 \rangle + \langle h_0^2 \rangle = \lambda \langle h_0^2 \rangle - \frac{E_{ext}}{A} + \langle h_{ext}^2 \rangle. \tag{5}
\]

From equation (5), recalling equation (2) and the equipartition theorem [30], we can express the out-of-plane rippling ratio between the cases with and without external excitation,
where \( k \) is the kinetic energy they carry. The magnitude of the externally induced ripples and the amount of dynamic rippling in the form of a competition between the strategies the effect of external excitation on the overall amount of flattening is possible only in a non-adiabatic case, when \( 1 + 2E_{\text{ext}}/k_B T_0 \), the effect of flattening cannot be observed, because \( \varepsilon = \sqrt{1 + (\langle h_{\text{ext}} \rangle / \langle h_{\text{0}} \rangle)^2} > 1 \), guaranteed for a non-zero excitation amplitude, as expected. Note that flattening may still be possible even if the energy is conserved. With strong anharmonicity, \( \langle h_{\text{ext}} \rangle \neq 0 \) must be present in equation (3). In addition, the values of \( A \) may no longer be identical between equations (2) and (3) and thus flattening may occur as a result of anharmonic stabilization of the lattice long-range order [9, 7].

For the harmonic system we are considering, however, flattening is possible only in a non-adiabatic case, when \( 1 < \lambda < (1 + 2E_{\text{ext}}/k_B T_0) \), according to equation (6). The value of \( \lambda \) is not freely adjusted to fit the simulation data; its value is obtained directly from the simulations. It is important to keep in mind that \( \lambda \) depends on the external excitation, as well as the membrane’s thermal conductivity (i.e. graphene specific) and the dissipative properties of the heat sink represented by the thermostat. Ultimately, the externally induced portion \( \langle h_{\text{ext}} \rangle \) of the overall vibrations must be obtained in terms of the source frequency \( \omega \) and the amplitude \( F_0 \); the external contribution to the kinetic energy is then automatically obtained, because \( E_{\text{ext}} \propto \langle h_{\text{ext}} \rangle \omega^2 \).

We obtained the following closed-form approximation for the MSD of the externally excited ripples:

\[
\langle h_{\text{ext}} \rangle \simeq \frac{1}{4} \frac{F_0^2}{m^2} (\gamma^2 \omega^2 + (\omega^2 - \omega_0^2)^2),
\]

where \( \gamma \) is the effective damping in graphene as a result of energy loss in the membrane, \( f_0 = \omega_0/2\pi = 0.29 \) THz is the resonant response frequency of the graphene membrane with dimensions as stated earlier, and \( m \) is the mass of the carbon atom. Above, \( \omega \) is the frequency of external excitation. The kinetic energy carried by these ripples is \( E_{\text{ext}} = \frac{m}{2} \langle h_{\text{ext}} \rangle \omega^2 \). The derivation and the approximations made are described in section 1 of the supplementary information (available at stacks.iop.org/Nano/24/055701/mmedia). Note that \( \langle h_{\text{ext}} \rangle \) and \( E_{\text{ext}} \) oscillate at \( 2\omega \). Here, \( \gamma \) is a fitting parameter, because no direct thermostatting was applied to and thus no effective viscosity was preset for the simulated sample, as described earlier. Also, regardless of the boundary conditions (periodic or clamped), the presence of harmonic functions results in an expression similar to equation (7). Even if the membrane shape is different, the expression of interest is identical to equation (7), except for the pre-factor on the right different from 1/4.

Note that the frequency regions of interest (see figure 2(a)) decrease with the effective sample size as \( 1/L \) (see section 1 of the supplementary information available at stacks.iop.org/Nano/24/055701/mmedia). Thus, compared with the nanometer-sized samples, these regions are expected to shift toward giga- and megahertz for micron- and millimeter-sized samples, respectively, consistent with experiment [17]. The ripple magnitudes (both thermal and especially externally induced—see equation (7)) are also expected to increase with the increased sample size in the corresponding frequency regions of interest. Finally, one must also realize that the electronic contribution to the dynamic response to the excitation is not included in our simulations. Clearly, this contribution can be significant, especially at high excitation.

**Figure 3.** Simulated and analytically calculated normalized out-of-plane RMSD \( \varepsilon \) (left Y-axis) and simulated temperature (right Y-axis) (a); Fourier transforms of \( h_{\text{tot}}^2(t) \) at a constant excitation frequency and various amplitudes (b), and at a constant excitation amplitude and various frequencies (c).
 frequencies. We therefore believe that our results are qualitatively closer to those one may expect at significantly lower frequencies of excitation, applied to larger samples.

Shown in figure 3(a) is the normalized simulated RMSD of the out-of-plane displacement \( \varepsilon = \sqrt{\frac{h_{\text{out}}^2}{h_0^2}} \) at \( F_0 = 4 \) nN versus frequency from figure 2(a), alongside that obtained from equations (6) and (7). Each point in the analytical curve was obtained by using the simulated temperature shown in figure 3(a). In all cases, we found \( \lambda = T/T_0 < (1 + \frac{2F_0}{T_h T_0}) \).

There are some general differences between the simulated and the analytically calculated data, including the positions of the local minima. These differences are attributed mainly to the relative simplicity of our analytical model, compared to the atomistic simulation. Specifically, the response of the membrane is represented by the single (1, 1) mode (see section 1 of supplementary information available at stacks.iop.org/Nano/24/055701/mmedia) and no anharmonic effects are taken into account. Finally, no discrete atomistic effects (especially important for the highest values of \( \omega \) and thus short wavelengths) are taken into account. However, we see that flattening \( \varepsilon < 1 \) is confirmed by the analytical model, and the overall trends are qualitatively similar in the frequency range where flattening occurs.

The underlying cause for the observed flattening is revealed by examining equation (6), given equation (7) and strongly bearing in mind that \( E_{\text{ext}} \propto h_{\text{ext}}^2 \omega^2 \). Three main regimes of response to excitation are therefore identified: (i) low \( \omega \) and high \( h_{\text{ext}}^2 \); (ii) intermediate \( \omega \) and \( h_{\text{ext}}^2 \); and (iii) high \( \omega \) and very low \( h_{\text{ext}}^2 \). It is useful to consider these cases in the context of the spectra in figures 3(b) and (c), where we show the Fourier transforms of the simulated \( h_{\text{out}}(t) = \frac{1}{N}\sum N(z_i(t))^2 \) with and without excitation. In case (i), when a low-frequency excitation is applied, the large-amplitude out-of-plane vibrations carry a small kinetic energy contribution \( E_{\text{ext}} \), and thus the excited membrane is subject to almost no extra energy loss. At the same time, a large-induced peak \( h_{\text{ext}}^2 \) is added to the spectrum at twice the excitation frequency, as shown in figure 3(c) for the excitation at 2.5 THz. As a result, the overall amount of rippling increases significantly, compared to the case without external excitation. In case (ii) the dissipative effect plays the major role, because the externally contributed kinetic energy \( E_{\text{ext}} \) is now quite high, while \( h_{\text{ext}}^2 \) is moderate. As a result, the slow thermal vibrations at frequencies below \( \sim 7.5 \) THz mostly responsible for the ripples are effectively lowered due to dissipation. At the same time, a moderate externally induced \( h_{\text{ext}}^2 \) peak is added to the spectrum. This ‘spectral intrusion’ can be readily observed in figures 3(b) and (c) for the excitation at 10 THz. Finally, in case (iii) of high excitation frequencies both \( h_{\text{ext}}^2 \) and \( E_{\text{ext}} \) contribute negligible amounts, because they decrease as \( 1/\omega^4 \) and \( 1/\omega^2 \), respectively (see the spectra of \( h_{\text{ext}}^2(t) \) in figure S1 in the supplementary information available at stacks.iop.org/Nano/24/055701/mmedia). Note the large anharmonic peaks present for \( F_0 = 8 \) nN, decreasing at lower excitation amplitudes in figure 3(b). This observation supports our previous remarks and agrees with the general discussion of anharmonicity in thermally fluctuating membranes [9]. It is noteworthy that equation (6) is limited in determining the value of \( \varepsilon \) as a function of varying temperature. This limitation is mainly due to \( h_{\text{ext}}^2 \) given by equation (7), which is a function of the effective viscoelastic damping \( \gamma \), expected to depend on \( T_0 \). As a result, although we were able to use equation (6) to establish a general qualitative picture in figure 3(a) at \( T_0 = 300 \) K, the temperature dependence of \( \varepsilon \) may suffer, demonstrating the estimative nature of our analytical model. Consider, for instance, the results in figures 4(a) and (b), where we plot \( \varepsilon(T_0) \) obtained from simulation and from using equation (6), respectively. We used a constant excitation force of \( F_0 = 2 \) nN and plotted the temperature dependence for the excitation frequency of 1, 5, and 10 THz. A constant value of \( \gamma \) was used, identical to that used to obtain the results shown in figure 3(a). We can see that qualitative agreement between the simulation and equation (6) only exists at \( f = 1 \) THz, when, as explained earlier, \( h_{\text{ext}}^2 \) is large, while \( E_{\text{ext}} \propto h_{\text{ext}}^2 \omega^2 \) is negligible and thus causes little dissipation (analytically, the dissipative term in the denominator of equation (7) is small). As a result, equation (6) for low excitation frequency is approximated as \( \varepsilon \propto \sqrt{\langle h_{\text{tot}}^2 \rangle/T_0} \), and with \( \langle h_0^2 \rangle \propto T_0 \) we expect a decreasing trend with increasing temperature, reproduced in both figures 4(a) and (b). The situation changes when the excitation frequency increases: both \( E_{\text{ext}} \) and the damping-dependent term in equation (7) increase as \( \omega^2 \). Not surprisingly, the agreement between simulated data and equation (6) becomes poor. Even if we assume ideal dissipation in all cases (\( \lambda = 1 \)), the apparent discrepancies observed can still be attributed to the effective viscosity being assumed to be constant and independent of temperature. Although obtaining the \( \gamma(T_0) \) dependence is beyond the scope of this work and the preceding discussion is presented to demonstrate the limitations of equation (6), it also highlights the necessity of finding out graphene’s effective viscosity as a function of temperature.
In addition to flattening below the thermal limit, the case of low excitation frequency mentioned above is equally noteworthy. Shown in figure 4(a) are the frequency-dependent patterns obtained for the rectangular membrane. To further quantify these patterns, we calculated the reciprocal space spectra of $h^2_{\text{tot}}(x, y)$, similar to the frequency domain data presented in figures 3(b) and (c). Shown in figure 5(b) are the spectra normalized with respect to $\langle h^2_{\text{tot}} \rangle$ to bring out the relative strength of the externally excited ripples, compared to the overall amount of rippling. The curves in figure 5(b) have also been artificially shifted relative to each other along the $Y$-axis for further clarity. The observed peaks correspond to half the wavelength of the ripples (due to the square of the quantity of interest) and have been carefully compared with the ripple spacings, as calculated from the membrane structure snapshots in real space. The inset in figure 5(b) shows the ripple wavelength $\lambda_{\text{ripple}}$ as a function of excitation frequency. The dashed line in the inset is a rough fit using a simple $\lambda_{\text{ripple}} = c_{\text{eff}}/f$ relationship for an interference pattern, which is what is essentially observed, despite the thermal smearing. Here, $c_{\text{eff}} = 5207 \text{ m s}^{-1}$ is the effective speed of sound, reasonably close to the 6000 m s$^{-1}$ we used in our previous calculations (see section 1 of the supplementary information available at stacks.iop.org/Nano/24/055701/mmedia). We must note that a perfect interference pattern is only possible with ideal boundaries (ideally transparent in our case; also valid for an ideally reflecting boundary). The reality, however, is that the boundary is partially reflecting and transmitting due to the thermostat in the simulation. In an experiment, the same effect is expected from the energy loss due to substrate, depending on both the substrate itself and graphene–substrate interface properties. In addition, one can see that the externally excited peaks have significant width and thus represent a considerable departure from harmonic response, as also mentioned earlier in the discussion of the vibrational spectra in the frequency domain.

The effect of changing temperature goes beyond simple thermal ‘smearing’ of the externally driven ripples. Shown in figure 5(c) is a family of normalized spatial spectra obtained for various temperatures. The relative smearing of the externally excited ripples by the thermal oscillations is well expected, especially given the fact that the external excitation force amplitudes were not particularly high and the externally excited kinetic energy components correspond to only a fraction of the thermal component, as shown earlier in figure 3(a). Such thermal smearing is clearly observed as the temperature increases, with externally excited ripples significantly smeared at 400 K. However, the peaks associated with the externally excited ripples shift and split with temperature without a clearly defined trend, as shown in figure 5(c). The effective ripple wavelength range is 2–4 nm in response to an identical external excitation and is a result of temperature changes. This behavior cannot be explained from the standpoint of harmonic response and, in addition to the boundary losses (which change with temperature, as discussed earlier), is likely due to significant anharmonic coupling between thermal and driven ripples. A precise analytical relationship between temperature, the strength of anharmonic coupling and its overall effect on the spatial

\[
f = 1 \text{ THz}
\]
\[
f = 2.5 \text{ THz}
\]
\[
f = 3.75 \text{ THz}
\]
\[
f = 5 \text{ THz}
\]
\[
q (\text{Å}^{-1})
\]
\[
norm (\text{arb. units})
\]
distribution of the externally excited ripples presents solid ground for further work. For additional data on the spatial distribution of the ripples and their short-wavelength behavior from the normal–normal correlations, also see section 4 of the supplementary information (available at stacks.iop.org/Nano/24/055701/mmedia).

Note that these controllable patterns appear for both periodic and reflective finite boundaries (see figure S2 in the supplementary information available at stacks.iop.org/Nano/24/055701/mmedia). Also, we should note that the only requirement for obtaining these rippling patterns is that the region where the external force is applied is significantly smaller than the characteristic size of the membrane. The effective size of the excitation region will limit the shortest excitable ripple wavelength and may cause additional anharmonic effects (also see section 5 of the supplementary information available at stacks.iop.org/Nano/24/055701/mmedia).

In the case of atomically thin layers, surface morphology directly affects the electronic properties. Therefore, there exists a possibility for modulation of the electronic properties in terms of the reported dynamic patterns. Although the presence of a thermal component makes these patterns less regular than in an ideal interference pattern, they do preserve periodicity and amplitude, depending only on the source frequency, as demonstrated earlier. Such a possibility is extremely attractive, given the recent report of experimentally observed quantum confinement as a result of local strain-induced pseudomagnetic fields [18]. Because these fields are present both at the boundary and throughout the locally suspended layer [19], one could probe the effect of dynamic patterns on the quantum Hall effect in the suspended graphene sheets, depending on the external modulation. Moreover, as supported by the observation of strong nonlinear peaks in figure 3(b), generation of lattice solitons [31], or nearly isolated lattice vibration packets, should be possible. In addition to the possibility of dynamically controlled quantum dots via patterned lattice distortions, the optical response of the material may also be dynamically tunable. The local morphology of free-standing layers has a significant effect on the plasmon response of these materials [32–35] and therefore the described dynamic patterns may be utilized for control of their optical properties.

4. Conclusions

In summary, we have described control over the dynamic morphology of free-standing graphene membranes by applying a local sinusoidal out-of-plane excitation in the presence of thermal effects. Phenomena ranging from overall flattening well below thermal limit to stable dynamic ripple patterns, depending on the excitation amplitude and frequency, have been described. The external excitation could probably be applied to a region significantly smaller than the entire membrane optically (with a pulsed laser), mechanically, or via a local time-varying magnetic field to a current-carrying graphene sample. The use of such local morphology control may result in an effective modulation of the electrical and optical properties of atomically thin layers. Although our work was focused on graphene, the results are generally applicable to atomically thin layers in terms of using a time-varying force for studying their dynamic responses and modulating their properties.

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