

# Molecular dynamics simulation of synchronization in driven particles

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Synchronization plays a key role in many physical processes. We discuss a simple numerical model of synchronized particles using molecular dynamics simulations for particles moving through a viscous liquid. Modeling synchronization with simulations and table-top experiments can provide insight to complex behaviors in the natural world. We drive particles across a washboard potential energy landscape, a common technique to isolate complex synchronized patterns in both simulations and experiments. Our results show a variety of synchronization effects in single particle systems, which we characterize by plotting particle velocity as a function of applied driving force and phase diagrams of position versus velocity.

## I. INTRODUCTION

Synchronization is a universal phenomena in which individual oscillators change frequency due to external stimuli [1]. The flickering patterns of candle flames mediated by temperature fluctuations [3], vibrations of singing wineglasses interacting through sound waves [4], and metronomes vibrating through a supporting platform [5] are examples of in-phase coupled oscillations. Biological systems benefit from cooperative synchronization – birds coordinate wing flaps to optimize energy use during flight [6], frogs alternate croaking patterns [7], humans clap in time with music [8] and at a cellular level, neurons simultaneously fire in cardiac muscle [9] and brain tissue [10]. External forcing can cause or regulate synchronization – an electrical pacemaker regulates a heart beat and a pulsed light modifies the flashing pattern of fireflies .

Synchronized phase-locking or mode-locking first appeared in the scientific literature with Huygens’ 1665 experiments on motions of synchronized pendula in wall-mounted clocks [2]. A locked-mode is an integer frequency ratio. In Huygens clocks, the pendula were observed swing at same rate – a 1:1 mode. For simple pendula of different lengths, a 2:1 mode can be seen when a pendulum four times the length of another swings with twice the period.

Complex dynamics can be studied with colloid particles in experiments or simulations. Typical colloids are plastic spheres suspended in de-ionized water or silica beads suspended in organic solvent. Because colloids are large and move slowly particle position can be imaged in real time with an optical camera [15]. Experimental measurements of step-by-step dynamics of colloids performing phase-locked motions are useful for understanding synchronization at a single particle level [26].

Light is a tool for manipulating the colloidal environment. A circular optical trap uses radiation pressure from a laser beam to confine colloids [30]. A colloid centered in the beam is uniformly bombarded by photons while

off-center colloids collide with photons unevenly, creating a net force. Depending on the location of the particle in the trap, the radiation pressure either moves colloids toward the beam center or ejects it from the trap. An optically trapped colloid executing Brownian motion is a useful probe of microscopic forces [31]. Diffraction gratings can be used to create more complex light environments, such as periodic patterns of minima suitable for synchronization studies [29].

A single particle oscillator in a potential well is like a skateboarder in a half-pipe or a child on a swing. How a particle moves in an optical landscape can change when subject to external driving forces. A confined oscillator may synchronize its location to the periodic pattern of the external drive, moving back and forth in time with the beat. With a constant or dc drive the landscape modulates the particle velocity, below some threshold the dc force is not strong enough to push the particle across a potential maximum so the average velocity is zero, a phenomena referred to as pinning [24]. Above the pinning threshold, a particle subject to a constant drive force will increase its speed at a rate proportional to the external drive. A sinusoidal or ac drive can create mode-locking, where the average particle velocity is fixed for a range of dc drive forces [25].

Here we perform numerical studies on the synchronized dynamics of confined particles driven over a washboard shaped potential energy landscape. We chose this model for its relevance to condensed matter systems and ease of simulation. We describe our molecular dynamics model for a single particle in Section II. We summarize our results including synchronized motion of a single confined particle driven across a periodic landscape in Section III. We include exercises for interested students in Section IV. In Section V we describe how our results apply to physical systems such as dusty plasmas, superconducting vortices and Josephson junctions.

## II. MOLECULAR DYNAMICS SIMULATION

We use a classical model for studying the dynamics of  $N$  interacting particles, using the net force on each particle to calculate its trajectory. Particles are confined

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in a two-dimensional (2D) simulation of area  $A = L \times L$  where  $L = 36.5a_0$  where  $a_0$  is a dimensionless unit of length. An individual particle  $i$  has position  $\vec{r}_i = x_i\hat{x} + y_i\hat{y}$  and velocity  $\vec{v}_i = d\vec{r}_i/dt$ . The edges of the system are treated with periodic boundary conditions such that a particle leaving the edge of the system is mapped back to a position within the simulation boundaries by the transformation  $x_i + L \rightarrow x_i$  and  $y_i + L \rightarrow y_i$ . We show a schematic of the system in Fig. 1(a).

We confine the particles using a position dependent potential energy function, called a landscape or substrate. The landscape is modulated in the  $y$ -direction with the periodic function

$$U(y) = U_0 \cos(N_p \pi y / L) \quad (1)$$

where  $N_p$  are the number of periods, and  $U_0$  is an adjustable parameter to set the depth of the minima with simulation units of energy  $E_0 = 1$ . We plot this function in Fig. 1 for  $N_p = 3$ . In Fig. 1(a) we show the  $x-y$  plane with a contour plot of  $U(y)$  to illustrate the 2D potential energy landscape, where the maxima are colored red and the minima colored blue. The confining force on a particle  $i$  is calculated as  $\vec{F}_i^l(\vec{r}_i) = -\nabla U_l(\vec{r}_i)$ . In Fig. 1(b) we plot the function  $U(y)$  to illustrate how the magnitude  $\vec{F}_i^l$  is calculated from particle position  $y_i$ .

Particles are subject an external time-dependent driving force  $\vec{F}_D(t)$  applied parallel to the  $y$ -direction. We model this force as

$$\vec{F}_D(t) = [F^{dc} + F^{ac} \sin(\omega t)]\hat{y}, \quad (2)$$

with modifiable parameters including a constant component  $F^{dc}$ , and a time dependent component with amplitude  $F^{ac}$  and frequency  $\omega = 2\pi f$ .

The inertia of small particles is reduced by interactions with fluid particles [16]. In our model we assume colloids are overdamped – i.e. colloids are suspended in a continuous viscous fluid that dissipates energy such that the particles do not accelerate. Energy dissipation from the fluid is modeled with a frictional force on particle with velocity  $\vec{v}_i$  as  $\vec{F}_i^{drag} = -\eta\vec{v}_i$  with a friction coefficient  $\eta$  proportional to the fluid viscosity. We discuss friction models for spheres moving through fluids in Exercise IV A.

Newton's second law for an individual particle is simplified by the assumption  $\vec{a}_i$  is zero. The overdamped equation of motion for an isolated particle is

$$\eta\vec{v}_i = \vec{F}_i^l(\vec{r}_i) + \vec{F}_D(t). \quad (3)$$

with  $\eta = 1$  in units of  $v_0/F_0$ . The equation of motion provides a direct calculation of the velocity of an individual particle from its location  $\vec{r}_i$  and the simulation time.

The molecular dynamics simulation is controlled by a *for()* loop which runs from an initial to maximum integer number of time steps. Each integer time step interval

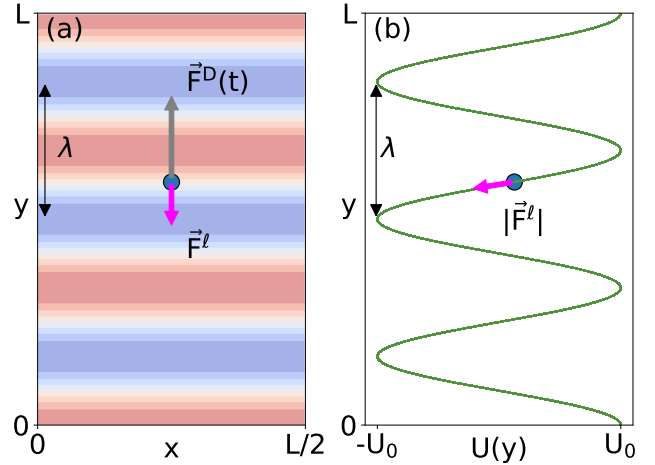


FIG. 1: Schematic of the simulation of a single particle driven across a washboard potential energy landscape. The period of the landscape is  $\lambda = L/N_p$ , where  $N_p = 3$ . The force due to the landscape is calculated from the gradient of the potential energy  $\vec{F}^l = -\nabla U(\vec{r})$ . (a) View of the  $x-y$  plane. The time-dependent applied driving force  $\vec{F}^D$  is parallel to the  $y$ -axis. The landscape is shown with a contour plot, with maxima in the potential energy marked in red and minima marked in blue. A particle is shown in the region between minima and maxima subject to competing forces of the landscape and applied driving force. (b) The potential energy function along the  $y$ -axis  $U_l(y)$ . The particle in (a) is shown at the same  $y$ -position. The slope of  $U_l(y)$  is the magnitude of force  $\vec{F}^l$ .

represents a simulation time of  $\Delta t$  simulation units  $\tau$ . At each time step we evaluate the net force on each particle as a function of its position  $\vec{r}_i(t)$  and then integrate the equation of motion to move particles to an updated position. Since the acceleration is zero, the integration of the equation of motion is performed via the Euler method

$$\vec{r}_i(t + \Delta t) = \vec{v}_i(t)\Delta t + \vec{r}_i(t) \quad (4)$$

for a small time step  $\Delta t = 0.001\tau$ . In Exercise IV C we describe the numerical methods for solving differential equations. The units of the simulated variables are summarized in Table I.

### III. MODE-LOCKING OF A SINGLE PARTICLE

A single particle responds to the applied driving force by moving across the landscape synchronized with the period of  $F^D(t)$ . The numerical implementation of the landscape is calculated with Eq. 1 as

$$F_y^l(y) = -A_p \sin(N_p \pi y / L) \quad (5)$$

where the force is scaled with parameter  $A_p = N_p \pi U_0 / L$ . In this section, we fix the landscape parameters to  $A_p = 0.1$  force units with  $N_p = 20$  minima corresponding to a spatial period  $\lambda = 1.825a_0$ .

TABLE I: Simulation parameters and units and comparable experimental values [26]. Blanks indicate units, rather than parameters that are not easily represented by a single value.

Quantity	Simulation Unit	Experimental value
length	$a_0 = 1$	$a_0 \sim 1.5\mu\text{m}$
energy	$E_0 = 1$	
force	$F_0 = E_0/a_0$	
time	$\tau = \eta a_0/F_0$	
velocity	$a_0/\tau$	$v \sim 5\mu\text{m/s}$
driving period	$T = 100\tau$	$T \sim 1.3\text{ s}$
viscosity coefficient	$\eta = F_0\tau/a_0$	$\eta \sim 10^{-3}\text{ Pa}\cdot\text{s}$
substrate period	$\lambda = 1.825a_0$	$\lambda = 3.5\mu\text{m}$
substrate amplitude	$U_0 = 0.058E_0$	$U_0 = 75k_B T \sim 2J$
temperature	$T_0 = 0^a$	$T \sim 300K$

<sup>a</sup>to learn more about the effects of temperature on this simulation see Ex. IV D where we study the limit  $U_0 \sim k_B T$ .

The driven particle adjusts its motion to traverse the periodic landscape with the same rhythm as the applied time-dependent force  $F^d(t)$ . When  $F^d(t) > A_p$ , a particle can overcome the barrier height of the landscape, and the particle hops between minima in the energy landscape. In Fig. 2(a) we plot  $F^d(t)$  as a function of time with constants  $F^{dc} = 0.1$ ,  $F^{ac} = 0.05$  and  $f = 0.01$  cycles per time unit  $\tau$ . The temporal period of the driving force is  $T = 1/f = 100\tau$ . We mark the maxima of  $F^{ac}(t)$  with the vertical dashed lines – i.e.  $\sin(\omega t) = 1$ . In Fig. 2(b) we show the  $y$ -position of the particle as a function of time. We normalize  $y$  by  $\lambda$  to show particle motion between substrate minima. The initial particle position is  $y = 0$ . As the driving force increases, the motion of the particle synchronizes with the driving period  $T$ . We measure average particle velocity  $\langle v_y \rangle$  as the average displacement over the period of the driving force. The particle moves in the positive  $y$ -direction through  $\Delta y = 2\lambda$  over the time period  $T$ , with the average velocity  $\langle v_y \rangle = 2\lambda f$ . The horizontal dashed lines coincide with the condition the particle is at a substrate minima where substrate force  $F^l = -A_p$ , i.e.  $\sin(\pi y/\lambda) = 1$ . The particle synchronizes its motion such that the driving force is maximum when the landscape force is minimum, as indicated with the intersection of vertical and horizontal lines. The slope of  $y/\lambda$  is proportional to the net force on the particle. This can be seen in Fig. 2 where local extrema appear in  $y/\lambda$  when  $F_D(t) \approx 0$ .

The competition between the driving force and landscape potential can produce a variety of hopping patterns in the particle motion. The relative values of  $F^{ac}$ ,  $F^{dc}$  and  $A_p$  control the rate and distance a particle moves forward and backward in the landscape. To explore the possible hopping patterns, we sweep through a range of  $F^{dc}$  for fixed  $F^{ac}$  and  $A_p$  to illustrate distinct oscillation modes. A mode is a periodic pattern of hops with

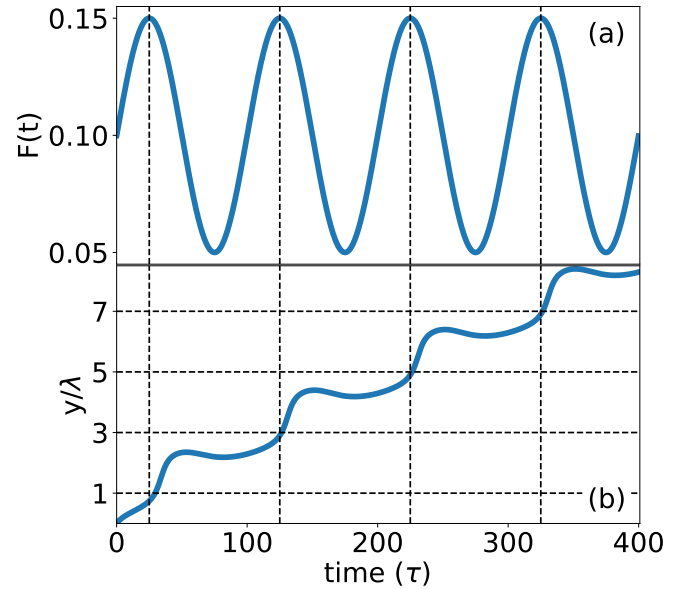


FIG. 2: (a) The applied driving force  $F^d(t)$ . (b) The  $y$ -position of the driven particle normalized by the period of the substrate  $\lambda$ .

a constant average particle velocity,  $\langle v_y \rangle$  over a range of driving forces  $F^{dc}$ . In Fig. 3 we keep  $F^{dc}$  constant for  $10^5$  simulation timesteps and measure the average velocity  $\langle v_y \rangle$  as a function of  $F^{dc}$ , with  $F^{ac} = 0.05$ ,  $f = 0.01$ , and the landscape shown in Fig. 2. The modes appear as discrete steps in  $\langle v_y \rangle$  across a range of  $F^{dc}$ ,

Each step represents a different pattern of hops between substrate minima performed by the particle due to the landscape confinement. In Fig. 3 at low  $F^{dc}$  the average velocity  $\langle v_y \rangle$  is zero. Since  $A_p$  is large compared to the extrema of  $F^d(t)$ , the particle oscillates back and forth in a single minima with no net velocity. At higher  $F^{dc}$  the particle velocity  $\langle v_y \rangle$  increases in steps of uniform height,  $\langle v_y \rangle = n\lambda f$ , where  $n$  is an integer. The step width is non-linear and can have a variety of interesting patterns such as a devil's staircase related to chaotic dynamics [12]. We note the individual steps are not strictly horizontal since particles move continuously through the landscape, so changes in  $F^{dc}$  will create linear changes in net particle displacement that are observable within a single step.

We explore a range of parameters in Ex. IV E where we perform this sweep for several values of  $F^{ac}$  including the limiting case that the applied force is a constant value. When  $F^{ac} = 0$ , no steps appear in the  $\langle v_y \rangle$  vs  $F^{dc}$  curve. When  $F^{dc} - F^{ac} < -A_p$  the particle moves in the negative  $y$ -direction during part of the cycle. The period of the substrate can be used to control the intrinsic velocity of the dc driven particle, and the applied ac drive can cause mode-locking which appears as non-linear steps in the force-velocity relationship.

In the following section we include exercises to explore our numerical model, including the to linear drag equa-

tion in Ex. IV A, the equation of motion in Ex. IV B, and numerical integration techniques in Ex. IV C. An extended model that includes finite temperature effects is in Ex. IV D. We explore changes in the particle motion with parameter changes in Ex. IV E. Methods to visualize the dynamics of the particle motion are in Ex. IV G.

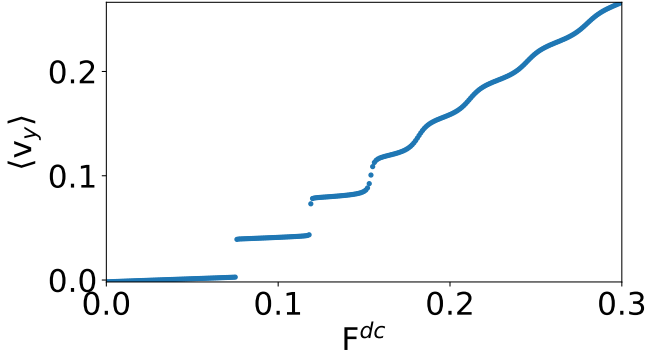


FIG. 3: Average particle velocity  $\langle v_y \rangle$  as a function of  $F^{dc}$ , the constant parameter of  $F^d(t)$  defined in Eq. 2. The remaining parameters in  $F^d(t)$  are as in Fig. 2,  $F^{ac} = 0.05$  and  $f = 0.01$ .

## IV. ASSOCIATED PROBLEMS

### A. Drag Models and Reynolds Numbers.

Stokes' law describes the viscous drag on a sphere moving at velocity  $\vec{v}$  as  $\vec{F}^{lin} = -3\pi\eta D\vec{v}$  where  $\eta$  is the dynamic fluid viscosity and  $D$  is the particle diameter [37]. In simulation we subsume the constants  $3\pi D$  such that  $3\pi D\eta \rightarrow \eta$ . Often drag forces are modeled as a polynomial series [37]

$$\vec{F}^{drag} = -b\vec{v} - cv^2\hat{v} + \dots \quad (6)$$

Truncating the series to the first term is justified by demonstrating the sphere has a low Reynolds number  $R = Dv\rho/\eta$  where  $\rho$  is the fluid density and  $v$  the particle's speed. When  $R$  is small, the quadratic and higher order terms may be ignored in favor of the linear drag term.

**Using reasonable values for the experimental analog of this system, show that the Reynolds number is small.** In addition to the values listed in Table I, the liquid density  $\rho$  varies with temperature with  $\rho \sim 10^3 \text{ kg/cm}^3$  a reasonable first order approximation [40].

### B. Equation of Motion

Newton's second law states that the acceleration of a particle  $i$  is proportional to the sum of forces on the

particle

$$m_i \vec{a}_i = \sum \vec{F}_i \quad (7)$$

where the constant of proportionality is the inertial mass  $m_i$ . The addition of a dissipative force to a dynamical equation of colloid motion is typically modeled with a drag force proportional to the particle's velocity in the opposite direction of motion  $\vec{F}^{drag} = -\gamma\vec{v}_i$  where  $\gamma = 3\pi\eta D$  is the drag coefficient described in Ex. IV A. The ratio of  $m/\gamma$  is known as the momentum relaxation time, and is small for particles with low Reynolds numbers. The mass of the typical experimental colloid particle is 15 picograms, leading to a momentum relaxation time on the order of microseconds. **Confirm for the values listed in Table I, the momentum relaxation time is  $m/\gamma \approx 0.5 \mu\text{s}$ .**

When  $m/\gamma$  is small, particle acceleration can be ignored entirely. (a) Using Newton's Second Law for a small momentum relaxation time, **show that a particle confined to a landscape exerting force  $F^l(\vec{r}_i)$  subject to a time dependent drive  $F^d(t)$  can be modeled with the equation of motion described in Eq. 3.**

### C. Integration Methods

To calculate the position of the particle we integrate the equation of motion using the standard definition of velocity  $\vec{v}_i = d\vec{r}_i/dt$  via the Euler method. Our equation of motion provides a direct calculation for particle velocity, as demonstrated in Ex. IV B. When solving ordinary differential equations, the Euler method is effective for solving linear equations of the form  $dy/dt = f(t, y(t))$  with initial condition  $y(t_0) = y_0$ . The solution is calculated algorithmically by stepping in time through  $n$  integer steps  $t_n = t_0 + n\Delta t$ . At each subsequent step the new value for  $y$  is calculated as a map solution using discrete times  $y_{n+1} = y_n + f(t_n, y_n)$ . **Apply the Euler method to our equation of motion to solve for the analytic expression of position of a particle  $y_n$  at the  $n$ th timestep (i.e. Eq. 4).**

The Euler method can be applied to non-linear equations effectively if the time step  $\Delta t$  is kept sufficiently small [38]. In our simulations we use the timestep  $\Delta t = 0.001$ , but we find for a single particle we can just as easily use timesteps  $\Delta t = 0.01$  and  $\Delta t = 0.1$  over the same maximum time (i.e. decrease the overall simulation time and sampling rate) and reproduce the motions shown in Fig. 2(b). This is due to the linear nature of the solution for an isolated particle. **Demonstrate the simulation timestep can be large with no change to the overall physics in the motion of a single particle.** You will need to modify the maximum simulation time, and the frequency of data saved for plotting purposes.

Interaction with neighboring particles makes a small timestep essential for molecular dynamics simulations.

Since particle-particle interactions are typically non-linear, the interaction force changes significantly over small distances, making this simulation a unique exception to typical simulation timesteps.

Often molecular dynamics algorithms are solved with higher order methods such as the Verlet or Runge-Kutta methods. These methods include higher order terms that provide accurate solutions for second order differential equations, i.e.  $dy/dt = v$  and  $dv/dt = f_2(y, t)$ . In Newton's second law  $f_2$  would be proportional to the net force on a single particle. In our model we do not integrate  $f_2(y, t)$  since  $a = dv/dt = 0$ , as described in Ex. IV B.

#### D. Brownian motion

Brownian motion is a phenomena in which visible particles change direction, apparently at random, due to collisions with invisible fluid particles. The rate of collisions depends on the temperature, viscosity and density of the suspending fluid. At higher temperatures the increased kinetic energy of particles makes collisions more likely, as described by the Maxwell-Boltzmann distribution [41].

In molecular dynamics simulations it is common to treat the invisible fluid particles as a continuous substance to reduce computational expense. Temperature effects can be modeled by applying randomized forces  $f^T$  to the visible particles. We use the normal distribution from the NumPy random module to generate a series of  $f_n^T$  values for each integer timestep  $n$  [42]. A normalized random distribution of forces causes fluctuations in motion equally in all directions such that the force  $f^T$  averaged over a finite time interval is zero. In one dimension this is expressed as

$$\langle f^T(t) \rangle = \frac{1}{N} \sum_n^N f_n^T = 0 \quad (8)$$

such that  $t = N\Delta t$ . Random normal distributions are not correlated over the independent variable, for the independent variable to be time

$$\langle f^T(t) f^T(\tau) \rangle = 2\eta k_B T \delta(t - \tau) \quad (9)$$

for times  $t$  and  $\tau$ , where  $\delta(t - \tau)$  is the Dirac-delta function. Here the scaling factor energy of  $k_B T$  is derived from the Boltzmann constant  $k_B$  and temperature  $T$  [36]. A particle with sufficient energy  $k_B T_{min}$  may hop over landscape barriers. In our simulations, we define temperature as  $k_B T/E_0 \rightarrow T$  with constants set to unity to compare directly with force.

**With no applied driving force, find the minimum temperature required for a single particle to hop over maxima in the potential landscape.** Assume the particle is confined to move along the  $y$ -direction, include Brownian motion in the model. In Fig. 4, we demonstrate the transition from pinned to hopping in a non-driven colloid for  $T/A_p = 4.0, 5.0, 6.0$ . The crossover to a hopping regime occurs at  $T/A_p \approx 5.6$ .

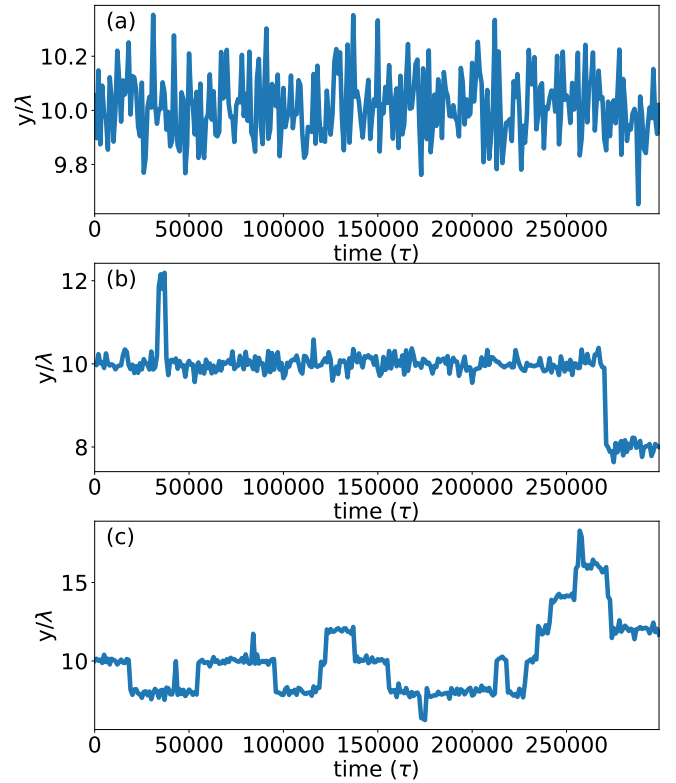


FIG. 4: Motion of a particle undergoing Brownian motion on a washboard potential energy landscape. (a)  $T/A_p = 4.0$  no hopping occurs. (b)  $T/A_p = 5.0$  a three hops occur across distances of  $\Delta y = 2\lambda$ . (c)  $T/A_p = 6.0$  many hops occur. Note the scale of the  $y$ -axis differs in each panel and the time scale is much greater than in Fig. 2.

In Fig. 4 we show the position versus time of a particle confined to a sinusoidal landscape undergoing Brownian motion. We increase the temperature relative to the amplitude of the landscape troughs until the particle enters a hopping regime. In Fig. 4(a) the temperature is  $T/A_p = 4.0$ . The particle executes a random walk centered at the potential minima  $y/\lambda = 10$ . We ran many simulations and never observed a hop to another minima within this simulation time. In Fig. 4(b)  $T/A_p = 5.0$  where hops between potential minima are possible but not probable, a single hop occurs from  $y/\lambda = 10$  to  $y/\lambda = 12$  at  $t \sim 40000\tau$  and  $y/\lambda = 10$  to  $y/\lambda = 8$  at  $t \sim 280000\tau$ . In other simulations with identical parameters, we sometimes observed no hops or two hops between minima, as expected with random fluctuating systems. In Fig. 4(c)  $T/A_p = 6.0$  we observe many hops. The random thermalized kicks are sufficiently large to make the particle perform what appears to be a random walk of hops atop the substrate.

For a driven particle, we ignore the effects of temperature in these simulations. We note that even at  $T/A_p = 6.0$  the hopping rate is much less than the frequency of the applied drive. This can be seen in the time



scale in Fig. 4, where the average hopping rate is approximately every  $15000\tau$ , a value much larger than  $T = 100\tau$ . At sufficiently high temperatures, Brownian motion does affect the formation of Shapiro steps and can be observed in experiments.

### E. Exploring Parameters of the Model

A range of oscillation behaviors can be explored by varying the relative strength of the confining landscape and external driving force. For the single driven particle, the hopping patterns are typically characterized by  $n_f$  the number of steps forward versus  $n_b$  the number of steps backward within a single time period. The total displacement of  $(n_f - n_b)\lambda$  is the net hop length. In Fig. 2, the particle moves forward through two minima ( $n_f = 2$ ), and does not move backward a full minima ( $n_b = 0$ ). In order to achieve backwards hops, the driving parameters must have a ratio of  $F^{ac}/F^{dc} > 1$  and a difference  $|F^{ac} - F^{dc}| > A_p$ . The frequency must be sufficiently low so that the applied force is large over a sustained time interval, allowing a particle to hop a substrate maxima. Experiments with single colloids [26] show hopping patterns with backward steps are more difficult to achieve than forward, with  $n_f = 0$  to 6 for  $n_b = 0$ ,  $n_f = 0$  to 5 for  $n_b = -1$ , and  $n_f = 2$  to 3 for  $n_b = -2$  were realized for a variety of applied forces parameters  $F^{dc}$  and  $F^{ac}$ . **Find the parameters that create hopping patterns similar to those in Ref. [26].**

In Fig. 5 we show the effects of lowering the frequency of applied driving force, holding the remaining parameters fixed to  $F^{dc} = 0.1$ ,  $F^{ac} = 0.05$ , and  $A_p = 0.1$ . We observe no backward steps ( $n_b = 0$ ) across a broad range of frequencies. The number of forward steps  $n_f$  increases as frequency decreases. Low frequencies result in sustained positive or negative motion of a particle, where in Fig. 5(e) the particle reaches the top edge of the system and wraps back to  $y = 0$ .

In Fig. 6 we fix  $f = 0.01$ ,  $F^{dc} = 0.1$ , and  $A_p = 0.1$ . We show the results for increasing  $F^{ac} = 0.2, 0.3, 0.4$ . Here  $n_b$  increases from  $n_b = 0$  in (a) to  $n_b = 2$  in (b) to  $n_b = 4$  in (c).

### F. Phase plots

Synchronization can be visualized with phase plots or Lissajous figures, first reported in 1857 [13]. A mode-locked system is confined to a closed loop in parametric phase space with patterns determined by the mode. (As we will demonstrate in phase plots in Figs. X, Y, Z).

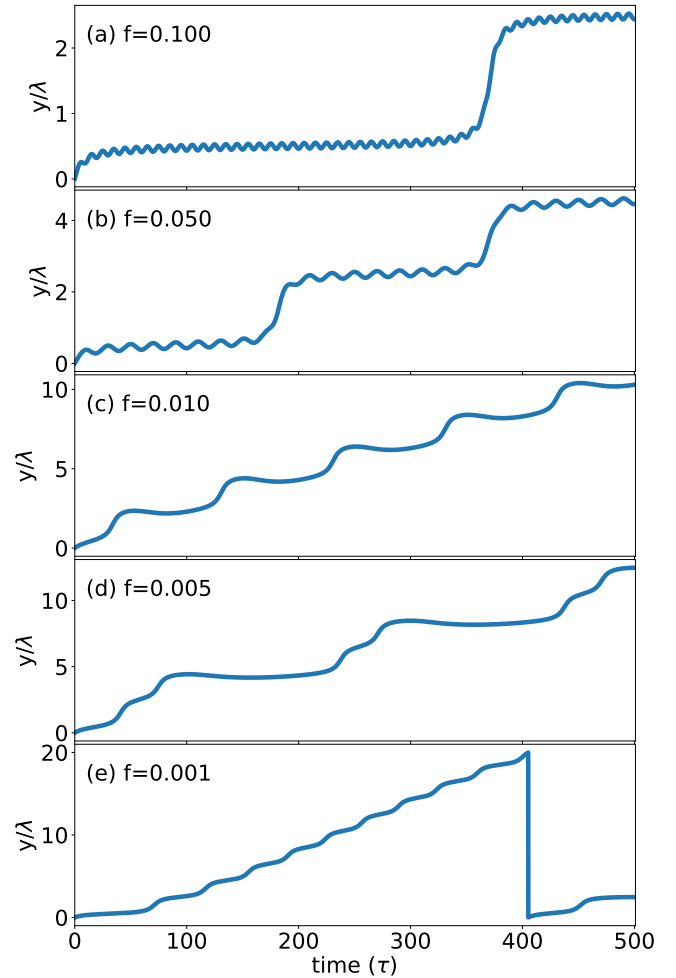


FIG. 5: Motion of a particle with the same values of  $F^{dc} = 0.1$  and  $F^{ac} = 0.05$  as Fig. 2 and decreasing frequency as marked. In (e) the particle wraps the periodic boundary conditions.

Note the term phase has two scientific meanings: oscillation phase  $\phi_0$  is the constant of the argument in a periodic function  $\sin(\omega t + \phi_0)$  and phase space is used to characterize oscillating systems with parametric plots of  $v_y$  vs.  $y$ .

### G. Visualizing motion

#### Drawing the contour

The code for generating a two dimensional colored plot of the potential landscape is calculated by evaluating the analytic function in Eq. 1 for a grid of values  $(x_n, y_n)$ .

*I have python code that does all of this, but I don't have it set up to play nice with my new python MD code*

Also the animation library in python...

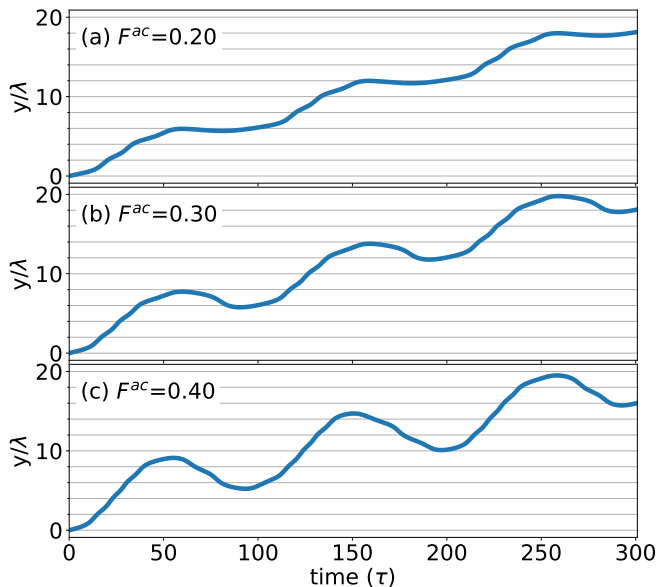


FIG. 6: Motion of a particle with the same values of  $F^{dc}$  and frequency as Fig. 2 and increasing  $F^{ac}$ . (a) The particle moves forward six minima with no steps backward. (b) The particle moves forward eight minima with two steps backward. (c) The particle moves forward nine minima with four steps backward.

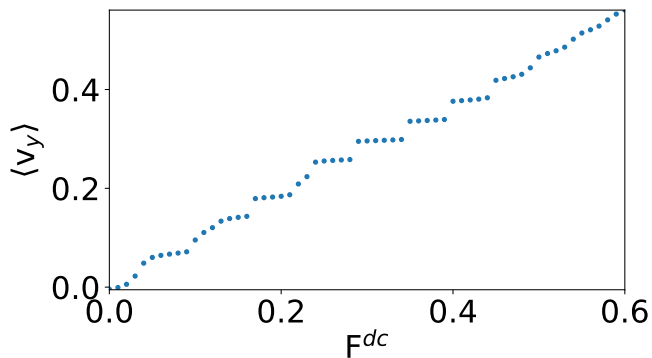


FIG. 7: Mode-locked steps in  $\langle v_y \rangle$  vs.  $F^{dc}$  of a particle with frequency  $f = 0.01$  and  $F^{ac} = 0.4$ . At low values of  $F^{dc}$  the particle may move backwards, distorting the steps.

## V. CONCLUSION

A single particle driven across a periodic potential landscape synchronizes its motion to environmental and external forces, a phenomena known as mode-locking. Our simulations reproduce experiments presented in Juniper *et al.* [26] of mode locking in driven colloids on a periodic optical landscape. Colloids are relatively easy to manipulate and image in experiments, making ideal proxies for systems such as cold atoms or electron gases [29]. Dynamical mode-locking is observed in quantum electronic devices as stepped regions in current-voltage (I-V) relationship, where voltage is the analog of external driving force and current is that of particle velocity. Known as Shapiro steps, these mode-locked or phase-locked currents have been observed due to applied ac voltages in single Josephson junctions [20, 22] and coupled arrays of junctions [21]. Shapiro steps vary in width depending on the strength of the applied ac forces, and are observed in a variety of ac and dc driven systems displaying non-Ohmic behavior in voltage-current curves, including charge waves, spin density waves and superconducting vortices in landscapes engineered with periodic patterns of pinning sites. Mode-locking is a useful probe of complex quantum mechanical systems since the motions of individual particles can only be inferred from other measurements. Our results can be relevant to synchronization effects in a broad range of experiments systems including optically confined colloids, superconductors with periodic pinning arrays, and the charge and spin of atomic systems.

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