

Chiral Swarmalators

Yichen Lu

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Contents

1	Models	2
1.1	Definitions	2
1.1.1	Self-propelled dynamics	2
1.1.2	Phase coupling dynamics	2
1.1.3	Order Parameters	2
2	Results	4
2.1	Single-chiral Swarmalators	4
2.1.1	Spatial Clustering in Circling State	5
3	Hydrodynamic Theory	5

1 Models

1.1 Definitions

1.1.1 Self-propelled dynamics

$$\dot{x}_i = v \cos \theta_i , \quad (1a)$$

$$\dot{y}_i = v \sin \theta_i , \quad (1b)$$

1.1.2 Phase coupling dynamics

- Additive coupling:

$$\dot{\theta}_i = \omega_i + K \sum_{j=1}^N f(r_{ij}) \sin(\theta_j - \theta_i + \alpha) , \quad (2)$$

- Mean-field coupling by oscillator number:

$$\dot{\theta}_i = \omega_i + \frac{K}{N} \sum_{j=1}^N f(r_{ij}) \sin(\theta_j - \theta_i + \alpha) , \quad (3)$$

which is similar to the swarmalator model.

- Mean-field coupling by oscillator density:

$$\dot{\theta}_i = \omega_i + \frac{KL^2}{\pi d_0^2 N} \sum_{j=1}^N f(r_{ij}) \sin(\theta_j - \theta_i + \alpha) , \quad (4)$$

where d_0 is the radius of the interaction circle.

Here, $f(r_{ij})$ is a function of $r = |\mathbf{r}_i - \mathbf{r}_j|$, and K is the coupling strength. The function $f(r)$ can be defined as

1. $f_H(r) = H(d_0 - r)$, $r_0 > 0$;
2. $f_E(r) = e^{-\frac{r}{d_0}}$, $r_0 > 0$.

The natural frequencies ω_i are distributed with following two cases:

1. **Single-chiral swarmalators:** The natural frequencies ω_i are distributed in $U(\omega_{\min}, \omega_{\max})$ for all swarmalators and $\omega_{\min} \omega_{\max} > 0$.
2. **Double-chiral swarmalators:** The frequencies are distributed in two symmetric uniform distributions, representing two types of chirality. Exactly half of the swarmalators have natural frequencies $\omega_i \sim U(\omega_{\min}, \omega_{\max})$ and the other half have natural frequencies $\omega_i \sim U(-\omega_{\max}, -\omega_{\min})$.

1.1.3 Order Parameters

Some order parameters can be introduced to measure the level of spatiotemporal coordinations among swarmalators and distinguish the different collective states of the system. Firstly, the usual order parameter to measure the global phase synchronization among swarmalators can be defined as the following complex function:

$$Z(t) = R(t) e^{i\psi(t)} = \frac{1}{N} \sum_{j=1}^N e^{i\theta_j(t)} , \quad (5)$$

where $i = \sqrt{-1}$. The degree modulus $R(t) = |Z(t)|$ is the absolute value of the mean of the complex numbers $e^{i\theta_i}$, which can be interpreted as the absolute value of the mean normalized velocity of all

swarmalators. When $R \approx 1$, swarmalators are fully phase synchronized, and when $R \approx 0$, swarmalators are phase incoherent.

The order parameter R is not enough to measure the emergence of partial clustered phase synchronization of swarmalators. Therefore a local order parameter can be introduced to measure the clustered synchronization:

$$Z_c^k(t) = R_c^k(t) e^{i\psi_c^k(t)} = \frac{1}{|C_k(t)|} \sum_{j \in C_k(t)} e^{i\theta_j(t)}, \quad (6)$$

N_c is the number of clusters, C_k is the k -th cluster, and $|C_k|$ is the number of swarmalators in the k -th cluster (see the details of the determination of clusters in Appendix). An averaged global order parameter can be further introduced as

$$R_c(t) = \frac{1}{N_c(t)} \sum_{k=1}^{N_c(t)} R_c^k(t). \quad (7)$$

As the swarmalators within the k -th cluster are fully synchronized, $R_c^k \approx 1$, and therefore also globally $R_c \approx 1$.

Except the study for phase coherence, another order parameter can be defined to further describe the locking of the frequencies of swarmalators under the chiral condition:

$$\Delta\Omega = \frac{1}{N_c} \sum_{k=1}^{N_c} \left[\frac{1}{|C_k|^2} \sum_{i,j \in C_k} \left(\langle \dot{\theta}_i \rangle - \langle \dot{\theta}_j \rangle \right)^2 \right], \quad (8)$$

where $\langle \dot{\theta}_i \rangle$ is the average phase velocity of the i -th cluster, which can be defined by

$$\langle \dot{\theta}_i \rangle = \lim_{T \rightarrow \infty} \frac{1}{T} \int_{t_0}^{t_0+T} \dot{\theta}_i(t) dt. \quad (9)$$

When $\Delta\Omega \approx 0$, the swarmalators will organize to form several clusters, within each cluster swarmalators are frequency-locked. Note that frequency locking is different from phase locking, as the frequency-locked swarmalators can have different phase velocities, which also means that they can have different rotational radii. Obviously, frequency locking is a more loose condition than phase locking.

2 Results

2.1 Single-chiral Swarmalators

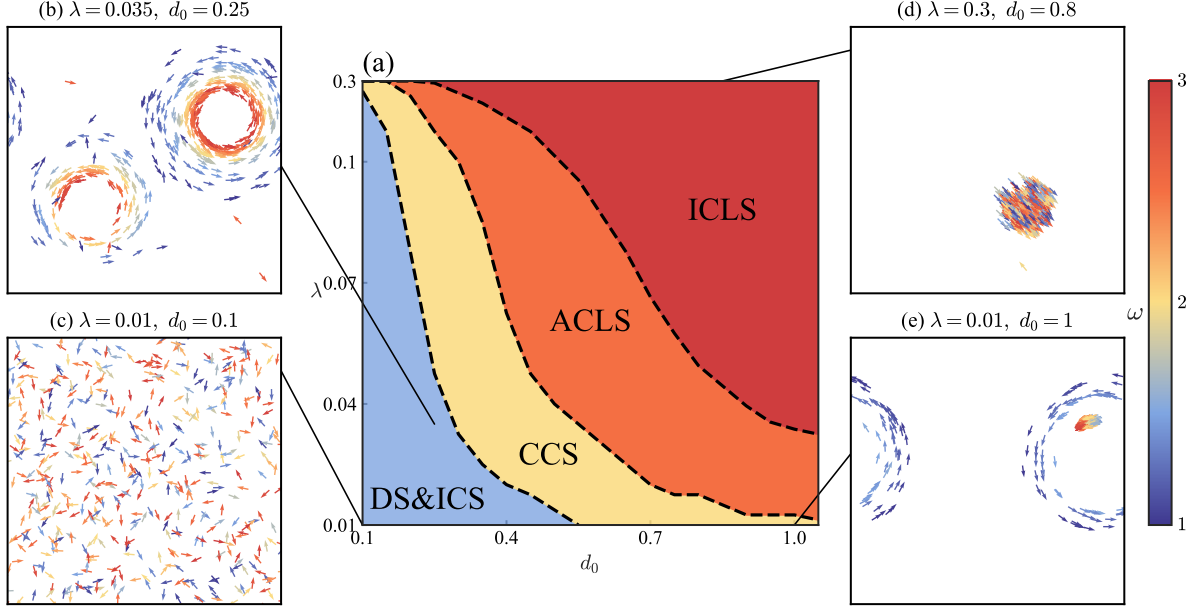


Figure 1: Phase diagram and snapshots of mono-chiral swarmalators. **(a)** Phase diagram in the (d_0, λ) plane. For the sake of clarity, the scale of λ and d_0 is non-uniform. **(b)**, The snapshots of CS ($\lambda = 0.08, d_0 = 0.1$). **(c)**, DS ($\lambda = 0.01, d_0 = 0.1$). **(d)**, CLS ($\lambda = 0.3, d_0 = 1$). **(e)**, CLS ($\lambda = 0.02, d_0 = 2$). The position and direction of each arrow are the instantaneous spatial position and phase of a swarmalator, respectively, and the color of the arrow denotes the corresponding natural frequency.

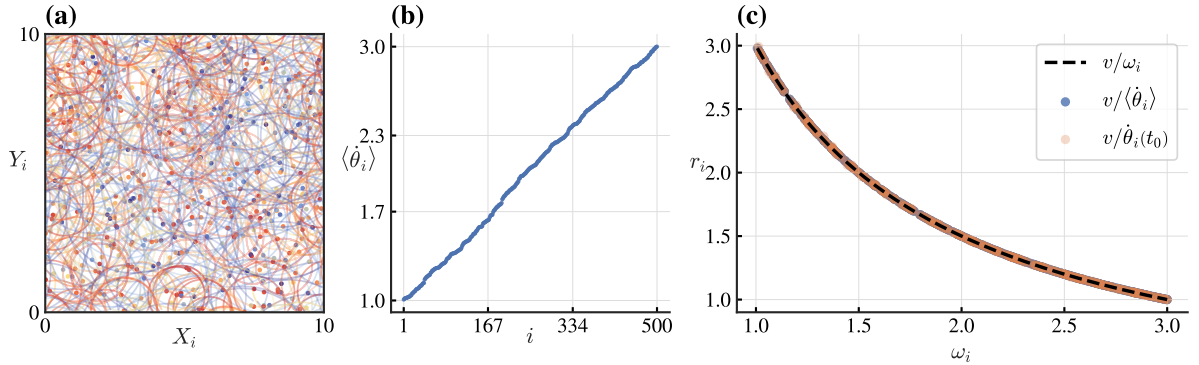


Figure 2: Swarming properties of DS for $\lambda = 0.01, d_0 = 0.1$. **(a)**: The orbits and the instantaneous rotation centers $\{c_i(t)\}$ of swarmalators. The color denotes the natural frequency. **(b)**: The profile of the averaged frequencies $\langle \dot{\theta}_i \rangle$. **(c)**: The instantaneous rotation radii r_i^{eff} and the average rotation radii $v/\langle \dot{\theta}_i \rangle$ against the natural frequencies ω_i . The black dashed line is the relation $r_i^0 = v/\omega_i$ for the uncoupled case.

2.1.1 Spatial Clustering in Circling State

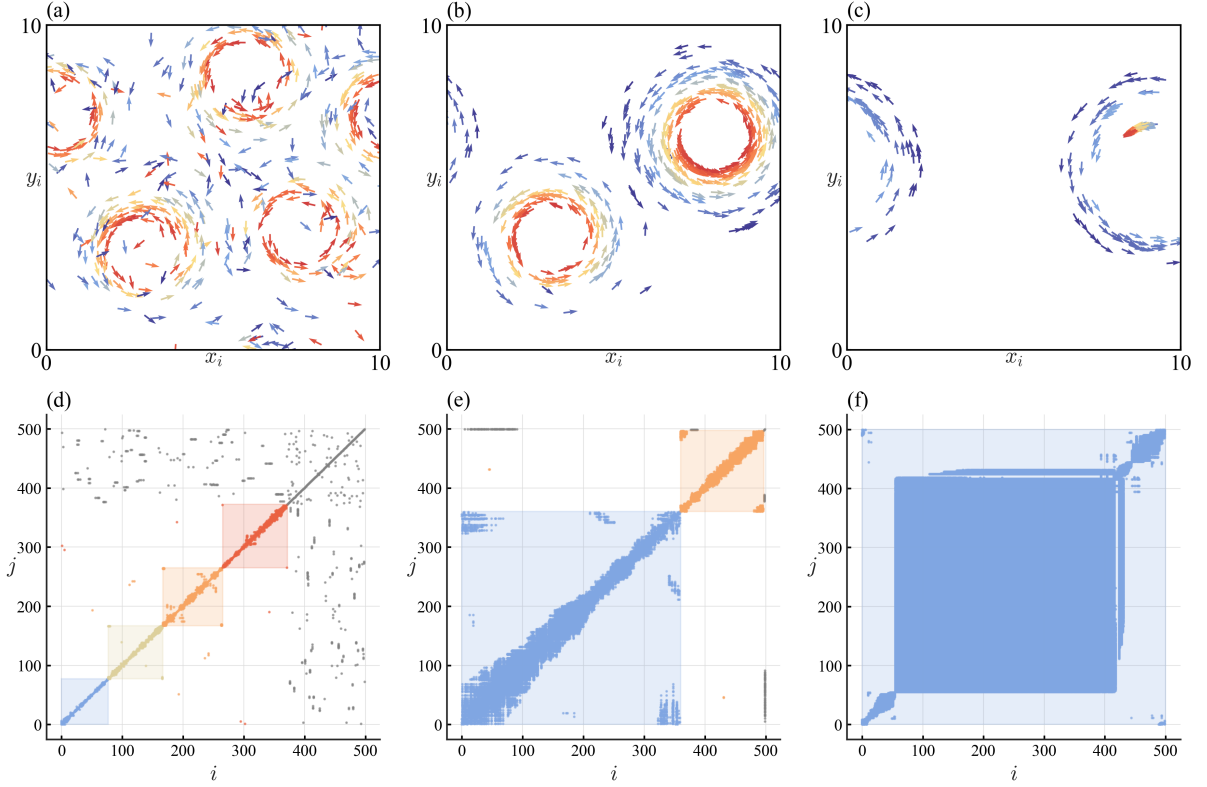


Figure 3: Spatial clustering in the circling state of single-chiral swarmalators. Top row: snapshots of the spatial distribution of swarmalators. Bottom row: the adjacency matrix A_{ij} of the network. **(a, d)**: $(\lambda, d_0) = (0.01, 0.35)$. **(b, e)**: $(\lambda, d_0) = (0.01, 0.65)$. **(c, f)**: $(\lambda, d_0) = (0.01, 1.05)$. In **(d-f)**, only those elements $A_{ij} = 1$ are plotted. The color of the matrix elements denotes different clusters, the gray elements represent the connections of the drifting swarmalators, and the white area represents the non-interacting region.

Sorting the swarmalators within each cluster by their standardized spatial angles, which are defined as

$$\arctan\left(\frac{y_i - \bar{Y}_i}{x_i - \bar{X}_i}\right), \quad (10)$$

where

$$\begin{bmatrix} \bar{Y}_i \\ \bar{X}_i \end{bmatrix} = \frac{1}{|C_k|} \sum_{i \in C_k} \mathbf{c}_i. \quad (11)$$

3 Hydrodynamic Theory

We replace the finite range alignment interaction by a pseudopotential (δ -interaction) in the model:

$$\begin{aligned} \dot{\mathbf{r}}_i &= v \mathbf{p}_i \\ \dot{\theta}_i &= \Omega_i + \lambda \sum_{j \neq i} \delta(\mathbf{r}_j - \mathbf{r}_i) \sin(\theta_j - \theta_i) \end{aligned} \quad (12)$$

where $\mathbf{p}_i = (\cos \theta_i, \sin \theta_i)$. Assuming that we have M species, consisting of N_1, \dots, N_M particles with identical frequencies $\tilde{\Omega}_1, \dots, \tilde{\Omega}_M$ respectively, and that N_1, \dots, N_M are all macroscopic in an area element over which macroscopic quantities (density, polarization) vary significantly, allows us to derive a

continuum theory for the particle dynamics. The combined probability density to find a particle of given species j at position \mathbf{r} with angle θ at time t is given by

$$f^{(j)}(\mathbf{r}, \theta, t) = \sum_{i=1}^N \delta(\mathbf{r} - \mathbf{r}_i(t)) \delta(\theta - \theta_i(t)) \delta_{\Omega_i, \bar{\Omega}_j}. \quad (13)$$

A Boltzmann-like equation for the combined density $f^{(j)}$ can be derived:

$$\begin{aligned} \dot{f}^{(j)}(\mathbf{r}, \theta, t) = & -\text{Pe} \mathbf{p} \cdot [\nabla_{\mathbf{r}} f^{(j)}(\mathbf{r}, \theta, t)] - \Omega_j \partial_{\theta} f^{(j)}(\mathbf{r}, \theta, t) + \partial_{\theta}^2 f^{(j)}(\mathbf{r}, \theta, t) \\ & - g \partial_{\theta} \left[f^{(j)}(\mathbf{r}, \theta, t) \int d\theta' \sin(\theta' - \theta) \sum_{i=1}^M f^{(i)}(\mathbf{r}, \theta', t) \right] - \partial_{\theta} \sqrt{2f^{(j)}(\mathbf{r}, \theta, t)} \eta_j(\mathbf{r}, \theta, t) \end{aligned} \quad (14)$$

where η_j represents spatiotemporal white noise with zero mean and unit-variance (the subscript j denotes that the noise-realization of a given ensemble is individual for each species).

In the following, we focus on a mean-field description and neglect the (multiplicative) noise term in Eq. (14). Now transforming Eq. (14) to Fourier space yields a hierarchy of dynamical equations for the Fourier modes $f_k^{(j)}(\mathbf{r}, t) = \int d\theta f^{(j)}(\mathbf{r}, \theta, t) e^{ik\theta}$ and $2\pi f^{(j)}(\mathbf{r}, \theta, t) = \sum_{k=-\infty}^{\infty} f_k^{(j)}(\mathbf{r}, t) e^{-ik\theta}$, reading

$$\begin{aligned} \dot{f}_k^{(j)}(\mathbf{r}, t) = & -\frac{\text{Pe}}{2} \left[\partial_x (f_{k+1}^{(j)} + f_{k-1}^{(j)}) - i\partial_y (f_{k+1}^{(j)} - f_{k-1}^{(j)}) \right] \\ & + (ik\Omega_j - k^2) f_k^{(j)} + \frac{igk}{2\pi} \sum_{m=-\infty}^{\infty} f_{k-m}^{(j)} F_{-m} \sum_{i=1}^M f_m^{(i)} \end{aligned} \quad (15)$$

From here, some lengthy algebra leads us to the following closed set of equations of motion for the fields $\rho := \rho^{(j)}$; $\mathbf{w} := \mathbf{w}^{(j)}$; $\Omega := \Omega_j$

$$\begin{aligned} \dot{\rho} = & -\text{Pe} \nabla \cdot \mathbf{w} \\ \dot{\mathbf{w}} = & -\mathbf{w} + \sum_{i=1}^M \frac{g\rho}{2} \mathbf{w}^{(i)} - \frac{\text{Pe}}{2} \nabla \rho + \frac{\text{Pe}^2}{2b} \nabla^2 \mathbf{w} - \frac{g^2}{b} \mathbf{w} \left(\sum_{i=1}^M \mathbf{w}^{(i)} \right)^2 \\ & + \Omega \mathbf{w}_{\perp} + \frac{\text{Pe}^2 \Omega}{4b} \nabla^2 \mathbf{w}_{\perp} - \frac{g^2 \Omega \mathbf{w}_{\perp}}{2b} \left(\sum_{i=1}^M \mathbf{w}^{(i)} \right)^2 + \mathcal{O}(\nabla \mathbf{w}^2) \end{aligned} \quad (16)$$

where $b = 2(4 + \Omega^2)$, $\mathbf{w}_{\perp} = (-w_y, w_x)$ and $\mathcal{O}(\nabla \mathbf{w}^2)$ represents all terms involving both derivatives and nonlinear terms in \mathbf{w} of arbitrary species and are therefore irrelevant for both the set of possible uniform solutions and for the linear stability of unpolarized (disordered) phases and can therefore be neglected for our purposes.