

**Simulation of Flexible Stems and Canopy Flow with
Immersed Boundary Method**

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Dedication

To those who witness my growth over the years

Abstract

In this thesis, we develop and apply numerical algorithms for simulation the interaction between fluid and thin flexible structures with sharp and diffused interface immersed boundary (IB) methods, respectively. With the sharp interface IB method, we simulate the interaction between air–water two-phase flows and thin flexible structures. In this algorithm, we propose a new method to avoid the unphysical fluxes across the thin structures, which effectively mitigates the issue of flow spurious penetration. In the proposed method, an explicit pressure boundary condition on the immersed boundary is applied. A new scheme for categorizing fluid, solid, and forcing nodes in the IB method is developed for avoiding extra pressure interpolation at the velocity projection step of the fractional-step method. Benchmark cases are tested and good agreement with the results in the literature is obtained. With the diffused interface IB method, we perform large-eddy simulation to investigate the dynamics of flow and canopy motions and the energy transfer in turbulent canopy flows. Different from the traditional approach that models the canopy as a continuous medium with a drag coefficient prescribed *a priori*, the IB method together with a beam model can explicitly capture the dynamics of individual stems. The periodic waving motion of the canopy, a phenomenon known as *monami*, is resolved. The simulation cases cover a broad range of stem flexibilities from rigid stems to oscillatory stems to stems yielding to the flow.

For flexible canopies, as the stem flexibility reaches a high Cauchy number, the stem fluctuation amplitude decreases, and thus, the canopy behaves more like a rigid canopy, a phenomenon we call ‘high flexibility-induced rigidity’, which is used to explain the similarities of the flow features between rigid and highly flexible canopies. Analyses of the turbulent kinetic energy (TKE) budget are performed. Inside the canopy, the wake production term associated with the canopy-induced spatial inhomogeneity is more pronounced than the shear production term. Near the canopy top in the flexible canopy cases, the waving term associated with the canopy drag–flow velocity correlation can be as large as one-half of the shear production term. Spectral TKE budget analyses further reveal dominant effects at two characteristic scales: the *monami* scale associated with the coherent structures in the mixing layer and the wake scale associated with the interval between adjacent stems. The waving term converts mean kinetic energy into dispersive kinetic energy and TKE. For the TKE in flexible canopies, the waving term is found to play an important role in the inter-scale and wall-normal transport terms. Our LES data show that the spectral shortcut mechanism proposed by previous studies is caused by the waving term, not the inter-scale transport term. The underlying mechanism responsible for the waving term is further elucidated through quadrant analyses of the correlation between the drag and velocity fluctuations of flexible canopy stems.

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Chapter 1

Introduction

1.1 Background and motivations

Fluid–structure interaction (FSI) problems are ubiquitous in biological and engineering applications [8]. Many problems involve the interaction between flow and thin flexible structures, such as waves impinging onto the ship hull, energy harvest from membrane vibration, hemodynamic around heart valve, and wind–plant interaction problems [9]. In these situations, the structures usually undergo large deformation or have complex geometry, which make the simulation extremely challenging [8].

The traditional approach for simulating FSI problems is the Arbitrary Lagrangian-Eulerian (ALE) method, where the boundary of the fluid grid fits the geometry of the structure surface. For moving or deforming structures with complex geometries, the ALE method may be complex as the mesh of the fluid solver needs to be regenerated

at every time step. Furthermore, the ALE method is prone to fluid grid failure when the structure deformation is large [8, 10]. In recent years, the immersed boundary (IB) method has been widely used in the computation of FSI problems. In the IB method, the grid does not deform, while forces are applied around the structure surface to satisfy the boundary conditions [11]. The IB method can be categorized into the diffused interface IB method and sharp interface IB method, depending on where the force is imposed [8]. In the diffused interface IB method, the surface force calculated on the immersed boundary is distributed to several adjacent fluid grid nodes by a smoothing function [12, 13, 14, 15, 16, 17]. As such, the boundary is essentially smeared. In the sharp interface IB method, the IB forcing is applied only on the grid nodes immediately adjacent to the immersed boundary [18, 19, 20, 1, 21, 22]. The sharp interface IB method can incorporate the wall-layer model for high-Reynolds-number flows [23, 24].

In this study, we solve two different kinds of flow–thin-structure interaction problems by applying the sharp and diffused interface IB methods, respectively. With the sharp interface IB method, we solve the interaction among air, water, and thin structures whose thicknesses are not resolvable by the background flow grid. The conventional sharp interface IB method causes unphysical flow momentum flux across the structure surface, which can induce water penetrating the structure when the structure thickness is unresolvable by the flow grid. We solve this problem by exerting a pressure boundary condition on the structure surface. The analysis and validation of our improved sharp interface IB method forms the first part of this thesis. With the diffused interface IB

method, we model the aquatic canopy by an array of flexible stems and simulate their interaction with the turbulent canopy flow. The periodic waving motion of the canopy, a phenomenon known as *monami*, is resolved. A broad range of stem flexibilities is simulated and their influences on the flow features are studied. In particular, we analyze the spectral TKE budget and study the effect of the drag-flow velocity correlation on the inter-scale energy transfer. These contents form the second part of this thesis.

1.2 Structure of this work

This work has two parts. The first part, which has Chapter 2, reports an improvement on the traditional sharp interface IB method for solving the interaction among air, water, and flexible thin structures. The numerical algorithm is described and its improvement over the tradition method is validated by two 2D cases and one 3D case. Chapters 3–5 are the second part, where a numerical study on the turbulent canopy flows with various stem flexibilities is reported. Chapter 3 describes the numerical algorithm, introduces the case setup, and defined the symbols for the result discussion in the following chapters. The numerical algorithm includes a flow solver, a stem solver, and a diffused interface IB method that couples the fluid and stem dynamics. Benchmark cases are reported for validation. Chapter 4 provides an overview of the canopy flow, including instantaneous flow field and canopy deformation, velocity and Reynolds shear stress profiles, stem deformation and *monami*, and the conditional average of flow velocity and canopy drag. Chapter 5 reports our spectral analyses on the TKE and DKE

budgets, where the spectral budget equation is derived and the effect of each term is analyzed. The mechanism of spectral shortcut, which involves the direct energy transfer between two characteristic flow scales, is discussed. Finally, Chapter 6 summarizes the conclusions of both parts.

Chapter 2

Solving Interaction among Air,

Water, and Flexible Thin

Structures with Sharp Interface

IB Method

2.1 Introduction

The interaction among air, water, and flexible thin plates is a typical problem in many science and engineering applications. Examples include aquatic vegetation flows in shallow waters, surface water waves interacting with ship hulls, and components of wave energy converters. A powerful tool for investigating this problem is numerical

simulation, which can provide detailed descriptions of the flow field and reveal the underlying flow physics over a range of temporal and spatial scales. Developing an accurate and robust numerical solver is needed but challenging, because it is required to capture the complex interaction dynamics among air, water, and flexible structure.

In fluid–structure interaction (FSI) problems with air–water two-phase flows, it is important to precisely capture the geometry of the air–water interface [25]. There are various approaches for capturing the interface between two fluid phases, for example, the level-set (LS) method, volume-of-fluid (VOF) method, and coupled level-set and volume-of-fluid (CLSVOF) method. In the LS method [26, 27, 28], a signed distance function ϕ from the interface is defined at each grid point, and the interface is described by the isosurface of $\phi = 0$. The VOF method [29, 30, 31, 32] defines the volume fraction of the fluid in each grid cell. The calculation of the interface normal direction is precise and straightforward in the LS method, while the VOF approach can conserve the mass of each phase of fluid [30]. In the CLSVOF method, the LS function is used to determine the interface-normal direction, which is used in the calculation of the flux of the VOF function; the VOF function is used to correct the value of the LS function to ensure mass conservation [33]. In the present study on the simulation of the interaction among air, water, and flexible plates, the CLSVOF method is employed to capture the air–water interface.

The geometry of the structures varies in different applications, for which the numerical method can be significantly different. In the present study, we confine our

interest to flexible thin plates. We review below some studies capturing flexible plates in numerical research of FSI problems. [34] and [35] coupled a fluid solver for direct numerical simulation and a structure solver for thin-membrane dynamics to study the arbitrarily large deformation of thin flags in uniform flow. They explored the effects of flag mass, flag bending rigidity, flag spanwise length, and flow Reynolds number on the instability of flag dynamics, and verified their results with the linear analysis. [36] integrated the simulation of two-phase fluid flows and flexible structures into a consistent finite element method (FEM) framework, where deformed unstructured grids were applied in both the fluid and solid phases. The LS method was adopted for tracking the air–water interface. [4] applied a particle finite element method (PFEM), where the governing equations for water motion and structure deformation were integrated into a consistent form, and both phases were described by Lagrangian trackers with different physical properties. [2] used the smooth particle hydrodynamics (SPH) method, where the water and deformable structure are discretized with particles tracked in a Lagrangian approach. [37] also modeled the water and structure using a Lagrangian approach, while the element bending group (EBG) method instead of SPH was used for the structure simulation. The interaction between the fluid and structure was implicitly captured by the inter-particle force between the SPH fluid particles and EBG structure particles at the interface. The studies reviewed above used either deformable mesh or meshfree Lagrangian methods for the fluid phase. [38] simulated the interaction between

a multiphase flow and flexible structure using the IB method. A refined mesh was defined on the surface of the structure to calculate the force on the immersed boundary. All of the above studies resolved the structure thickness in the flow simulation using high grid resolution. In some other scenarios, however, the structures are too thin to be resolved by the grid for the fluid simulation, such that they can only be treated as a thickless discontinuity in the flow field [1, 39].

In this chapter, we develop a numerical framework for the simulation of the interactions between two-phase fluid flows and flexible plates with unresolvable thickness. We propose a new method of explicitly applying a pressure boundary condition on the immersed boundary to mitigate unphysical momentum flux across the structure and to solve the fluid leakage problem that appears in the conventional IB method. We also propose a new algorithm for classifying the fluid, forcing, and solid nodes in the IB method to avoid an additional pressure interpolation at the velocity projection step of the fractional-step time-advancement method.

2.2 Method description

Our partitioned two-phase incompressible fluid solver [40, 41, 42] and large-deformation thin-plate structure solver [43] are coupled by the kinematic and dynamic boundary conditions on the immersed boundary, using the IB method [20, 1, 21] and our proposed explicit pressure boundary condition on the structure surface.

The fluid solver has both the direct numerical simulations (DNS) and large eddy

simulations (LES) capabilities. The two fluid phases are integrated into one system, and they are solved based on same equations while only distinguished by their different fluid properties, e.g. density and viscosity. The governing equations are

$$\frac{\partial \mathbf{u}}{\partial t} + \mathbf{u} \cdot \nabla \mathbf{u} = \frac{1}{\rho(\phi)} [-\nabla p + \nabla \cdot (\mu(\phi) \mathbf{S}) + \rho(\phi) \mathbf{g} - \nabla \cdot \boldsymbol{\tau}_{sgs}] + \mathbf{f}, \quad (2.1)$$

$$\nabla \cdot \mathbf{u} = 0, \quad (2.2)$$

where $\mathbf{u}$ is the fluid velocity, $\rho(\phi)$ is the fluid density, p is the pressure, $\mu(\phi)$ is the fluid dynamic viscosity, $\mathbf{S} = (\nabla \mathbf{u} + \nabla \mathbf{u}^T) / 2$ is the strain rate tensor, $\mathbf{g}$ is the gravitational acceleration, and $\mathbf{f}$ is the force exerted by the structure and decided by the IB method. Note that the fluid density $\rho(\phi)$ and fluid dynamic viscosity $\mu(\phi)$ are phase-dependent. In LES, the subgrid-scale stress $\boldsymbol{\tau}_{sgs}$ is determined by the dynamic Smagorinsky model [44, 45, 46]; in DNS, $\boldsymbol{\tau}_{sgs} = 0$. Surface tension is neglected because the problems of interest all feature a large characteristic length scale. The two-fluid interface is captured by the CLSVOF method [33]. The grid for the fluid solver is a staggered Cartesian grid, where the three velocity components are located on cell surfaces, while other scalar quantities are located at cell centers. A second-order central difference method is utilized for spatial discretization, and an explicit second-order Runge–Kutta method is applied for time integration. For each substep of the Runge–Kutta method, the governing equations of the fluid motion, Eqs. (2.1) and (2.2), are solved by the

fractional-step method [47],

$$\frac{\tilde{\mathbf{u}}^{k+1} - \mathbf{u}^k}{\Delta t} = \alpha_k \mathbf{H}^k + \beta_k \left(\mathbf{H}^{k-1} + \mathbf{f}^{k-1} \right), \quad (2.3)$$

$$\frac{\hat{\mathbf{u}}^{k+1} - \tilde{\mathbf{u}}^{k+1}}{\Delta t} = \alpha_k \mathbf{f}^{k+1}, \quad (2.4)$$

$$\frac{\nabla \cdot \hat{\mathbf{u}}^{k+1}}{\Delta t} = \alpha_k \nabla \cdot \frac{1}{\rho^k} \nabla p^{k+1}, \quad (2.5)$$

$$\frac{\mathbf{u}^{k+1} - \hat{\mathbf{u}}^{k+1}}{\Delta t} = -\frac{\alpha_k}{\rho^k} \nabla p^{k+1}, \quad (2.6)$$

where $k = 0, 1$ is the substep of the Runge–Kutta scheme; $\hat{\mathbf{u}}$ and $\tilde{\mathbf{u}}$ represent the intermediate velocity at different substeps of the fractional-step method, and $\hat{\mathbf{u}}^{k+1} = I(\tilde{\mathbf{u}}^{k+1})$, where I is the interpolation operator; $\mathbf{H}^k = -\mathbf{u}^k \cdot \nabla \mathbf{u}^k + \mathbf{g} + \frac{1}{\rho^k} \nabla \cdot (\mu^k \mathbf{S}^k)$ is the summation of the advection term, gravity term, and viscous term; $\alpha_0 = 1$, $\beta_0 = 0$, and $\alpha_1 = \beta_1 = 0.5$ are constants. For $\mathbf{X} = \mathbf{u}$ or p , $\mathbf{X}^{-1} = 0$, $\mathbf{X}^0 = \mathbf{X}(t_n)$, and $\mathbf{X}^2 = \mathbf{X}(t_{n+1})$, where n is the time step.

The structure solver is based on the Kirchhoff–Love thin-shell model and is solved by the finite element method [43]. The most significant feature of the present flexible plate solver is that only the translational degrees of freedom are involved for the nodal variables, different from the conventional method where the rotational degrees of freedom are also needed [43]. Therefore, the present solver significantly reduces the computational cost and allows using finer mesh for improving accuracy. Also, for the case where the time step for the structure solver is smaller than that of the fluid solver, we have applied a time-split scheme, where each time step of the fluid solver corresponds

to several time steps of the structure solver. In our FSI cases, the external force is the hydrodynamic force, which includes the pressure and viscous shear stress on the plate. Note that as the thickness of the structure is too small to be resolved by the grid of the fluid solver, the Froude–Krylov force and buoyancy force cannot be resolved as well and should be added to the total external force as two separate terms [48].

The external-forcing IB method is used in the our multiphase flow FSI framework to capture the influence of solid structure on the fluid flow [20, 1, 21]. In this method, the Eulerian grid nodes of the fluid solver are categorized into fluid, solid, and forcing nodes, where the forcing nodes are adjacent to the immersed boundary and located in the fluid phase, and the fluid and solid nodes are the rest of the grid nodes located in the fluid and solid phases, respectively. For thickness-unresolvable structures, there is no solid node, and the forcing nodes are located on each of the two sides of the structure. In the node classification algorithm of [1], a pair of adjacent cell-center or cell-surface nodes are both classified as forcing nodes, if the segment connecting them crosses the structure. Figure 2.1(a) shows an example using the method of [1]. In this algorithm, some cell-surface forcing nodes, such as Node D in Fig. 2.1(a), are adjacent to a pair of cell-center nodes located on the different sides of the structure. A problem emerges when the velocity at such cell-surface nodes is updated by the pressure gradient at the projection step of the fluid solver, as the pressure on the different sides of the structure should not communicate with each other. This problem can be solved by the proposed new node classification algorithm. The cell-center nodes are first classified using the

method of [1], and the classification of the cell-surface nodes is determined using the following criterion: (1) A cell-surface node is classified as a fluid node, if it is adjacent to two cell-center fluid nodes; (2) a cell-surface node is classified as a forcing node, if it is adjacent to a cell-center fluid node and a cell-center forcing node, or if it is adjacent to two cell-center forcing nodes that are located on the same side of the structure; (3) a cell-surface node is classified as a discontinuity node, if this node is adjacent to two cell-center forcing nodes located on different sides of the structure. An example of our proposed node classification algorithm for the thickness-unresolvable structure is shown in Fig. 2.1(b), where the classifications of Nodes C and D are different from the method of [1]. It is seen that the envelope of the discontinuity nodes fits the geometry of the structure and forms the inner and outer boundary of the forcing nodes located on inner and outer sides of the structure, respectively. Therefore, this envelope also marks the discontinuity of the pressure field, and is thus necessary for the application of the explicit pressure boundary condition on both sides of the immersed boundary described below. With the application of the pressure boundary condition, the velocity on a discontinuity node is not updated by the pressure gradient at the projection step of the fluid solver to avoid any unphysical flux across the structure. The classical node classification algorithm cannot distinguish the discontinuity node and the forcing node defined in our algorithm, and thus cannot decide if the velocity at a forcing node should be modified by the pressure gradient without additional information.

In the proposed IB algorithm, the velocity at a cell-surface forcing node is interpolated from the fluid nodes and immersed boundary to satisfy the velocity boundary condition on the immersed boundary. The velocity interpolation profile is linear for resolved wall layer or nonlinear for modeled wall layer depending on the Reynolds number and grid resolution [23, 24, 49], and the linear velocity interpolation algorithm is similar to [20] and [21]. After the velocity at the forcing nodes is interpolated, the velocity at the discontinuity nodes is also interpolated in the similar manner, while the interpolation stencil are the forcing nodes instead of the fluid nodes. Then, for a grid cell whose center is a forcing node, if all of its surface nodes, which are forcing nodes or discontinuity nodes according to their definitions, are on the same side of the structure, then that grid cell is divergence-free, because the velocities at these cell-surface nodes are finally interpolated from the fluid nodes located in the divergence-free flow field on the same side of the structure, and linear interpolation conserves the divergence-free condition. Note that the interpolation stencil of the discontinuity node is the forcing nodes, whose interpolation stencil is the fluid nodes; therefore, the velocity at the discontinuity nodes is essentially obtained from the interpolation from the fluid nodes. This divergence-free property of the forcing node is critical for moving structures, because our fractional-step algorithm requires a divergence-free history of the fluid nodes, and a fluid node can be a forcing node in the previous timestep if a nearby structure moves away from it. The velocity interpolation on the discontinuity nodes ensures that a cell-center forcing node is always divergence-free before it becomes a fluid node.

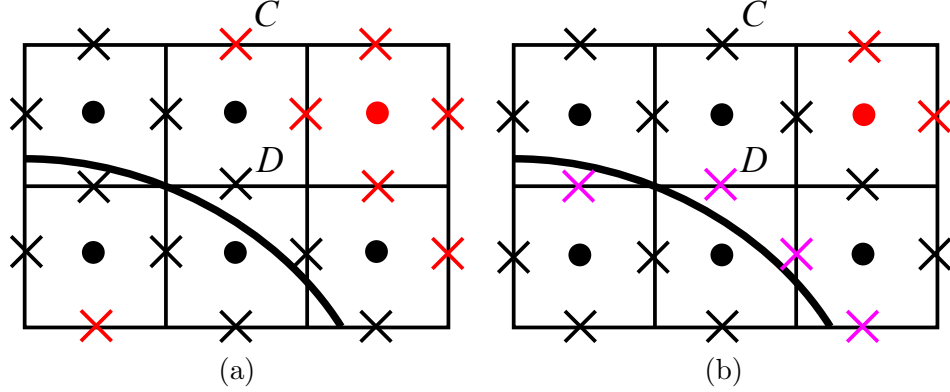


Figure 2.1: Node classification schemes for thickness-unresolvable structures in the immersed boundary method. (a) Scheme by [1]. (b) The present scheme. Thick curve represents the thickness-unresolvable structure, dot represents cell-center nodes, cross represents cell-surface nodes, red represents fluid nodes, black represents forcing nodes, and magenta represents discontinuity nodes. Note that Nodes C and D are classified into different categories between the previous and present schemes.

In the conventional external forcing IB approach, the pressure boundary condition is expected to be satisfied implicitly as a result of the correct velocity boundary condition. In such approach, however, the velocity defined on a cell-surface discontinuity node is corrected by the pressure gradient between the forcing nodes located on different sides of the structure, at the projection step of the fractional-step method as given by Eq. (2.6). Such a correction based on the pressure gradient across the thickness-unresolvable structure causes unphysical momentum flux. For example, for the case where a thickness-unresolvable plate confronting a flow stream, physically the pressure on the windward side of the plate is higher than that on the leeward side of the plate, such that there is pressure jump across the plate. In the fractional-step method, the pressure is solved by the Poisson equation (Eq. 2.5), and then at the projection step, Eq. (2.6), the pressure

jump across the plate induces a correction of velocity along the downstream direction at the discontinuity nodes, which induces an unphysical streamwise momentum flux across the plate. Similar issue also happens to the thickness-resolvable structures: [19] found unphysical momentum flux in their simulation of a cylinder crossflow, and this issue is caused by the unphysical pressure gradient across the surface of the cylinder. This leads to the following discussion: as the thickness-unresolvable structure is a discontinuity in the flow field, velocity and pressure boundary conditions should be applied on both sides of the structure; however, in the conventional external forcing IB approach, only the velocity boundary condition is applied by the velocity interpolation, while the significance of the pressure boundary condition is overlooked.

Therefore, we propose to solve this issue by explicitly applying the pressure boundary condition on both sides of the thickness-unresolvable structure. The pressure boundary condition is derived by projecting the momentum equation onto the immersed boundary,

$$\frac{D\mathbf{u}_\Gamma}{Dt} \cdot \mathbf{n} = -\frac{1}{\rho} \nabla p \cdot \mathbf{n} + \nu \Delta \mathbf{u} \cdot \mathbf{n} + \mathbf{g} \cdot \mathbf{n}, \quad (2.7)$$

where $\mathbf{u}_\Gamma$ is the velocity of the immersed boundary, $\mathbf{n}$ is the normal unit vector on the immersed boundary, and ν is the kinematic viscosity. Note that the viscous effects are often neglected [50, 20], and then the pressure boundary condition can be simplified to

$$\left. \frac{\partial p}{\partial n} \right|_\Gamma = -\rho \frac{D\mathbf{u}_\Gamma}{Dt} \cdot \mathbf{n} + \rho \mathbf{g} \cdot \mathbf{n}, \quad (2.8)$$

where the acceleration of the immersed boundary $D\mathbf{u}_\Gamma/Dt$ is evaluated by differentiating the velocity of the immersed boundary $\mathbf{u}_\Gamma$ with respect to time. Numerically, the pressure boundary condition is imposed on the cell-center forcing nodes. Similar to the velocity interpolation algorithm in the IB method, each cell-center forcing node is associated with an interpolation stencil that has one projection point on the immersed boundary and two/three fluid nodes for 2D/3D cases. It is assumed that the pressure distribution is linear near the immersed boundary, such that for 2D cases we have

$$p = a + bx + cy, \quad (2.9)$$

where a, b, c are coefficients. The pressure gradient along the wall normal direction is then expressed as

$$\left. \frac{\partial p}{\partial n} \right|_\Gamma = \nabla p \cdot \mathbf{n} = \frac{\partial p}{\partial x} n_x + \frac{\partial p}{\partial y} n_y = b n_x + c n_y, \quad (2.10)$$

where the wall normal unit vector (n_x, n_y) is defined at the projection point. Thus, we have

$$\begin{pmatrix} \partial p / \partial n \\ p_1 \\ p_2 \end{pmatrix} = \begin{pmatrix} 0 & n_x & n_y \\ 1 & x_1 & y_1 \\ 1 & x_2 & y_2 \end{pmatrix} \begin{pmatrix} a \\ b \\ c \end{pmatrix}, \quad (2.11)$$

where the subscript $i = 1, 2$ denotes the two fluid nodes. For a cell-center forcing node

of which the coordinate is (x_0, y_0) and pressure is p_0 , we have

$$p_0 = \begin{pmatrix} 1 & x_0 & y_0 \end{pmatrix} \begin{pmatrix} a \\ b \\ c \end{pmatrix} = \begin{pmatrix} 1 & x_0 & y_0 \end{pmatrix} \begin{pmatrix} 0 & n_x & n_y \\ 1 & x_1 & y_1 \\ 1 & x_2 & y_2 \end{pmatrix}^{-1} \begin{pmatrix} \partial p / \partial n \\ p_1 \\ p_2 \end{pmatrix}. \quad (2.12)$$

Equation (2.12) describes the pressure relation between the cell-center forcing nodes and the fluid nodes, and serves as the pressure boundary condition. This equation, along with the discretization form of the pressure Poisson equation for the fluid nodes, Eq. (2.5), can determine the pressure values of all the forcing and fluid nodes by solving a linear system. With this algorithm, the pressure boundary condition is explicitly implemented in the fractional-step method, and then in Eq. (2.6), the velocity at the cell-surface discontinuity nodes is corrected according to the pressure gradient determined by the pressure boundary condition instead of the unphysical pressure gradient across the structure. For the thickness-resolvable structures, as both the pressure and velocity boundary conditions are implemented on the immersed boundary, the solid nodes can be excluded from the computational domain, resulting in the decoupling of the solution from the fluid and the structure solvers [50, 51, 52].

The coupling algorithm for the flow and structure solvers is critical for the computational accuracy and numerical stability of partitioned systems and is dependent on the relative density between the fluid and structure [8]. When the density of the structure is much larger than that of the fluid, we adopt loose coupling, where both the fluid and

Table 2.1: Parameters of the case of collapsing water deforming a top-fixed flexible gate.

Variables	Values
Width of the water tank B	0.1 m
Initial water height in the tank H	0.14 m
Density of water ρ	10^3 kg/m^3
Kinematic viscosity of water ν	$10^{-6} \text{ m}^2/\text{s}$
Length of the rubber gate L	0.079 m
Thickness of the rubber gate t	0.005 m
Young's modulus of the rubber E	$4.0 \times 10^6 \text{ Pa}$
Poisson's ratio of the rubber ν_r	0.4
Density of the rubber ρ_{rubber}	$1.1 \times 10^3 \text{ kg/m}^3$
Gravity acceleration g	9.8 m/s^2

structure solvers update only once at each time step. This algorithm has low computational cost, but it becomes unstable when the density ratio between the structure and fluid is small (usually < 10) due to the added-mass effect [53]. In such situation, we switch to a strong coupling algorithm, where the fluid and structure solvers are iterated for converged solution at each time step [1]. All the test cases presented in Section 2.3 use the strong coupling because the density ratio between the structure and water is smaller than 10.

2.3 Test cases

2.3.1 Collapsing water column deforming a top-fixed flexible gate

This case was studied in the literature through experiments and simulations [2, 37, 54, 55]. As shown in Fig. 2.2, the water in the tank is initially blocked by a rubber gate with a fixed upper end and a free lower end. The parameters of this case are

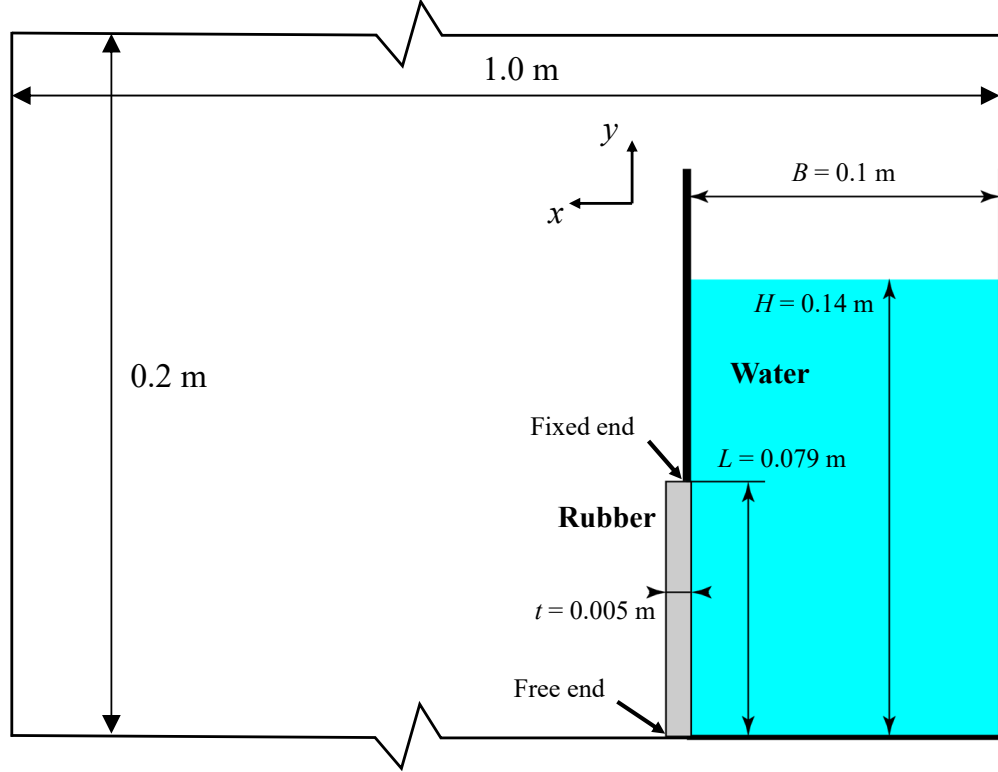


Figure 2.2: Case set-up of collapsing water column deforming a top-fixed flexible gate.

listed in Table 2.1. Both the water and rubber gate are initially stationary. The size of the computation domain is $1.0 \text{ m} \times 0.2 \text{ m}$, which is large enough such that the left boundary has negligible influence on the water flow for the time we are interested in. Free-slip conditions are applied to all the computational boundaries following [2]. The CFL number is 0.5. The length of the rubber gate is discretized with 40 Lagrangian grid cells, resulting in the resolution of the Lagrangian grid $\Delta s = 2.026 \times 10^{-3} \text{ m}$. The Eulerian grid for the flow is homogeneous inside and around the water tank, and stretched as the distance from the water tank increases.

Resolution independence test was performed on four different Eulerian grids. In these four cases, the grid resolution around the gate is $\Delta x = \Delta y = 1.563 \times 10^{-3}, 1.042 \times 10^{-3}, 0.781 \times 10^{-3}$ and 0.521×10^{-3} m, respectively. Figure 2.3(a) shows the time history of the horizontal displacement of the free end of the rubber gate obtained with different grid resolutions. It is seen that the curve for the lowest resolution shows some discrepancy with other curves when the displacement is large around $t = 0.13$ s, while the curves for the three higher resolutions agree well with each other, indicating that the results are independent of the grid resolution when the grid size is finer than 1.042×10^{-3} m. In the following content, the results of $\Delta x = \Delta y = 0.521 \times 10^{-3}$ m are demonstrated.

We compare the cases with and without applying the pressure boundary condition on the immersed boundary. Figures 2.4(c) and (d) compare the snapshot at $t = 0.32$ s between the two cases. It is seen from Fig. 2.4(c) that for the case without the pressure boundary condition, the water leaks from the plate. Also, because of the water leakage, the water level to the right of the gate is lower than that in Fig. 2.4(d), where the pressure boundary condition is applied. The time history of the horizontal displacement of the free end of the rubber gate for the two cases is compared in Fig. 2.3(b). It is seen that the pressure boundary condition influences the time history in two aspects. First, for $t < 0.15$ s, the plate deformation increases with time because of the pressure difference across the plate. For the case without applying the pressure boundary condition, the pressure difference across the plate is underestimated because of the water leakage,

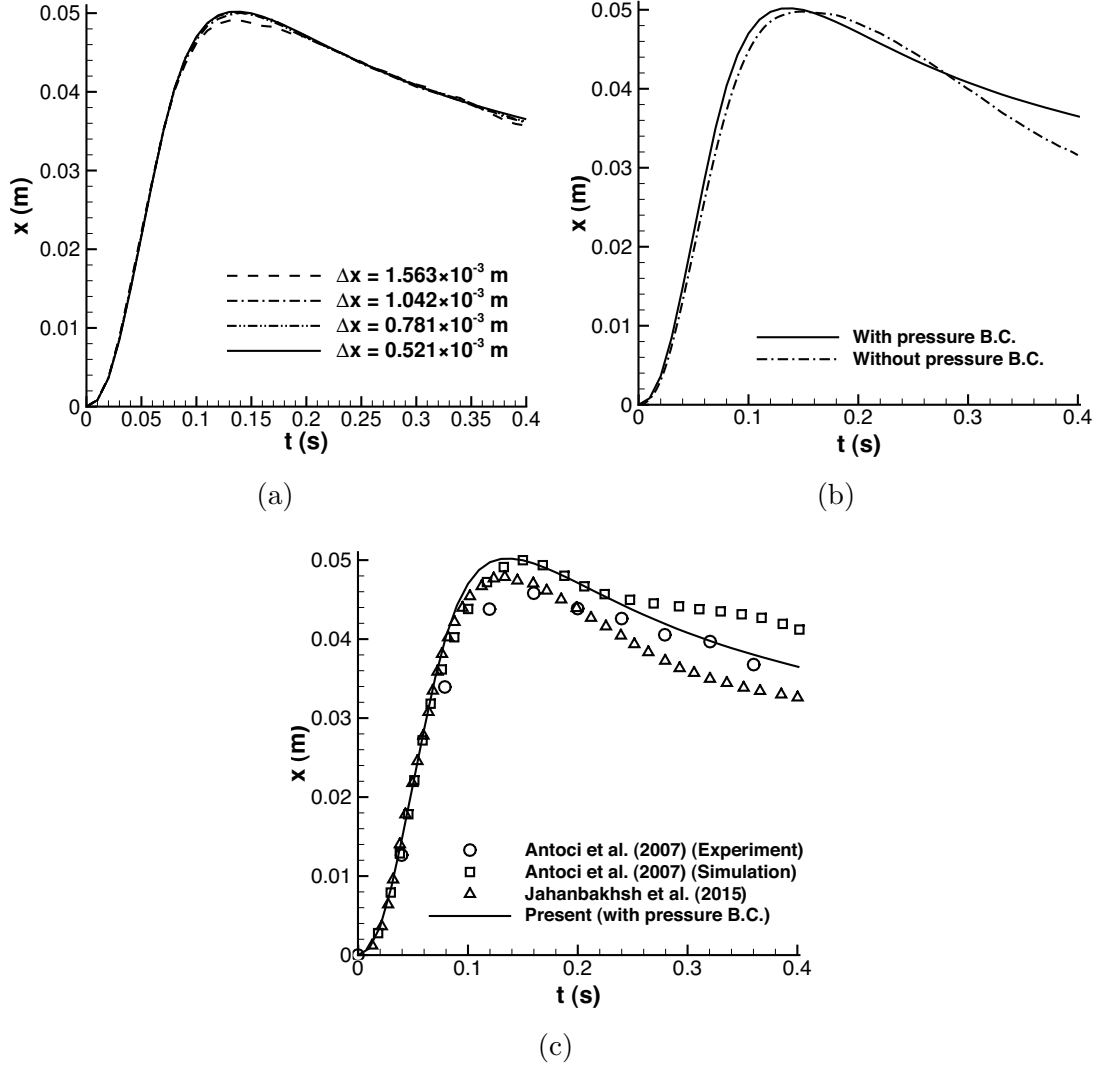


Figure 2.3: (a) Time history of the horizontal displacement of the free end of the rubber gate for different grid resolutions. (b) Time history of the horizontal displacement of the free end of the rubber gate for the cases with and without the pressure boundary condition applied on the immersed boundary. (c) Comparison of the time history of the horizontal displacement of the free end of the rubber gate between the present case and the experiment and simulation results of [2] and the simulation results of [3].

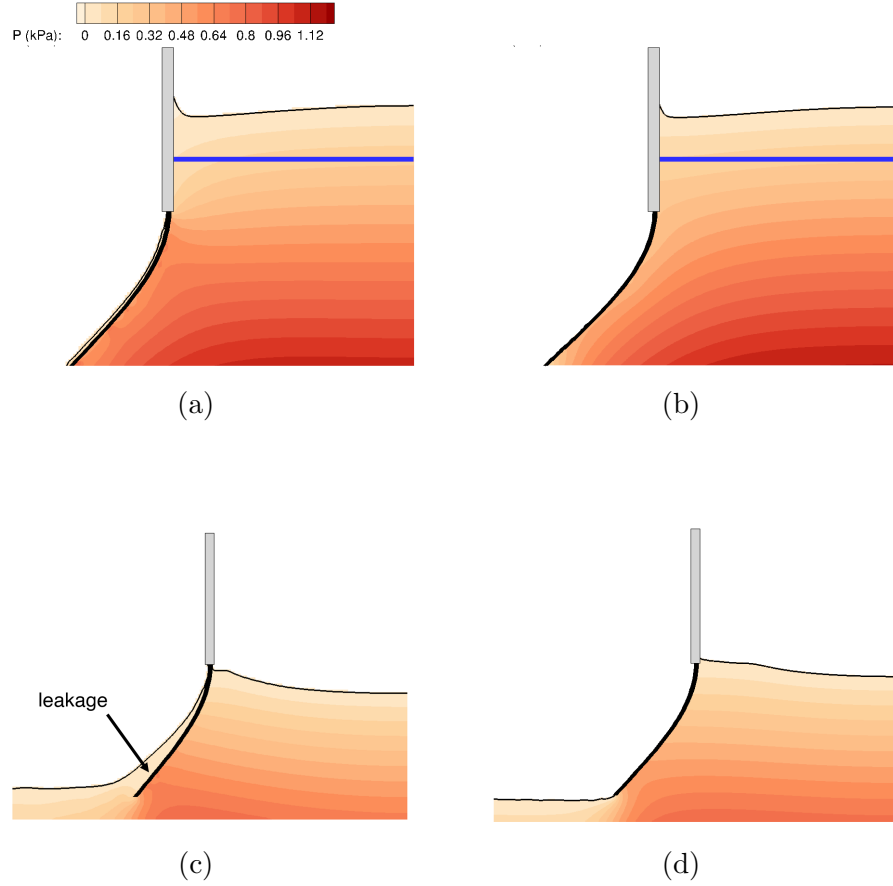


Figure 2.4: Snapshots for the simulation results of collapsing water deforming a gate. (a) The pressure boundary condition is not applied, $t = 0.08$ s. (b) The pressure boundary condition is applied, $t = 0.08$ s. (c) The pressure boundary condition is not applied, $t = 0.32$ s. (d) The pressure boundary condition is applied, $t = 0.32$ s. The blue lines in (a) and (b) show the location of $y = 0.1$ m.

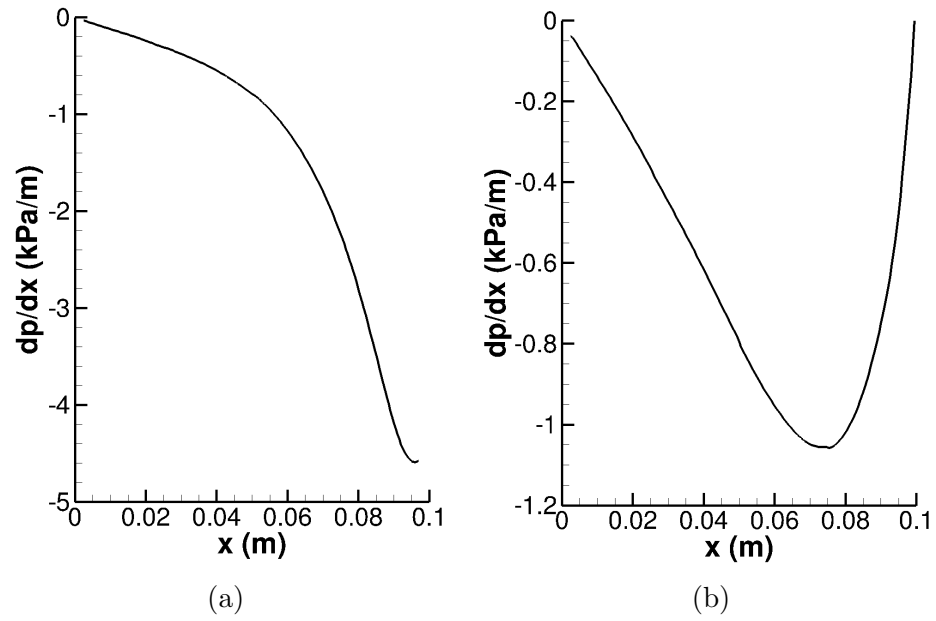


Figure 2.5: The pressure gradient along the line $y = 0.1$ m (shown as the blue lines in Fig. 2.4) when $t = 0.08$ s. (a) The pressure boundary condition is not applied, (b) the pressure boundary condition is applied.

so that the deformation rate is also underestimated. Second, for $t > 0.15$ s, the plate deformation decreases with time, because the resilience force of the plate becomes higher than the driving force provided by the pressure difference across the plate. For the case without applying the pressure boundary condition, because the pressure difference is underestimated due to the water leakage, the rebound rate of the plate is overestimated. As a result, the horizontal displacement in this time duration decreases faster in the case without applying the pressure boundary condition than in the case with the pressure boundary condition applied. Figures 2.5(a) and (b) show the pressure gradient along the line $y = 0.1$ m for $t = 0.08$ s in the cases without and with the pressure boundary condition, respectively. Note that the line touches the stationary part of the water tank wall at $x = 0.1$ m where the normal direction of the wall is $-\mathbf{e}_x$. Then, according to Eq. (2.8), $\partial p / \partial x$ should be zero at $x = 0.1$ m. It is seen from Fig. 2.5 that the pressure gradient at $x = 0.1$ m is noticeably different from zero in the case without the pressure boundary condition, while it is very close to zero in the case with the pressure boundary condition. From here, one can see that applying the pressure boundary condition is significant for obtaining the physical fluid flux across the structure boundary and the physical hydrodynamic pressure gradient on the structure surface.

The plate deformation and the shape of the free surface shown in Figs. 2.4(b) and (d) agree with the experiment result of [2] plotted in their Fig. 13. Figure 2.3(c) compares the time history of the horizontal displacement of the free end of the rubber gate obtained from the present simulation with the experimental and numerical results

Table 2.2: Parameters of the oil tank sloshing case.

Variables	Values
Domain length L_x	0.609 m
Domain height L_y	0.3445 m
Domain width L_z (for 3D case only)	0.039 m
Initial oil depth h_0	0.0574 m
Density of the oil ρ_l	917 kg/m ³
Kinematic viscosity of the oil ν_l	5×10^{-5} m ² /s
Length of the plate l	0.0574 m
Width of the plate w (for 3D case only)	0.0332 m
Thickness of the plate t	0.004 m
Young's modulus of the plate E	6.0×10^6 Pa
Poisson's ratio of the plate ν	0.49
Density of the plate ρ	1.1×10^3 kg/m ³
Gravity acceleration g	9.8 m/s ²

of [2] and the numerical result of [3]. Our result generally agrees with the previous experimental and numerical results.

2.3.2 Oil tank sloshing

This case was firstly studied experimentally and numerically by [4] and was later simulated by [56] and [57]. Figure 2.6(a) shows the initial state of the 3D simulation. The parameters for the case set-up are shown in Table 2.2. A fluid tank containing sunflower oil can rotate around a pivot axe along the z -direction going through the bottom-center of the tank. One edge of a flexible plate is clamped at the bottom-center of the tank. The length of the plate is the same as the initial oil depth. The numerical simulations in the literature were 2D. We perform both 2D and 3D simulations. The set-up of our 3D simulation follows the experiment of [4]. In our simulation, the origin of the coordinate is located at the center of the bottom wall. Free-slip conditions are applied

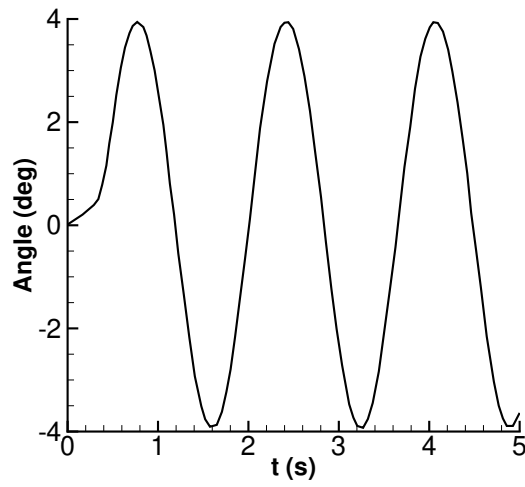
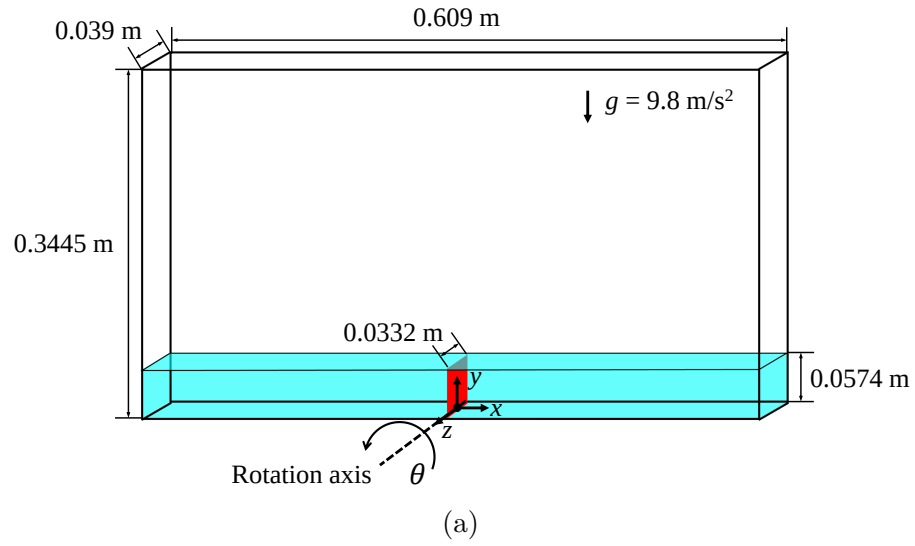


Figure 2.6: Set-up of the oil tank sloshing case. (a) Initial state of the 3D simulation; (b) time history of the rotation angle θ .

for all the boundaries following previous numerical simulations. The time history of the rotation angle in the experiment is shown in Fig. 2.6(b), where positive value indicates the counter-clockwise direction. The rotation angle varies sinusoidally with a period $T = 1.65$ s and amplitude $\theta_0 = 4^\circ$, except for the beginning of the experiment due to the transition from the rest state [4]. We fit the transition stage with piecewise polynomials. Our fitted formulas for the rotation angle are

$$\phi = \begin{cases} 1.301t & t < 0.2834 \text{ s}, \\ -1.352 + 26.61t - 147.0t^2 + 326.9t^3 - 225.1t^4 & 0.2834 \text{ s} \leq t \leq 0.5648 \text{ s}, \\ 4 \sin[2\pi(t - 2.015)/1.65] & t > 0.5648 \text{ s}, \end{cases} \quad (2.13)$$

where the unit of ϕ is degree and the units for all the coefficients are determined accordingly. In our simulation, the reference frame follows the tank rotation, such that the computational domain is fixed and the following two effects induced by the tank rotation are considered. The first effect is the inertia force. We added a volume force $\mathbf{f}$ on both the fluid and the structure to consider this effect,

$$\mathbf{f} = -\rho[\ddot{\boldsymbol{\phi}} \times \mathbf{r} + \dot{\boldsymbol{\phi}} \times (\dot{\boldsymbol{\phi}} \times \mathbf{r})], \quad (2.14)$$

where $\mathbf{r}$ is the position vector, and $\boldsymbol{\phi} = \phi \mathbf{e}_z$ with $\mathbf{e}_z$ being the unit vector along the spanwise direction. The second effect is the change in the gravity direction. The Eulerian mesh for the fluid domain is uniform near the plate with the grid resolution $\Delta x = \Delta y =$

Table 2.3: Grid convergence test for the oil tank sloshing case.

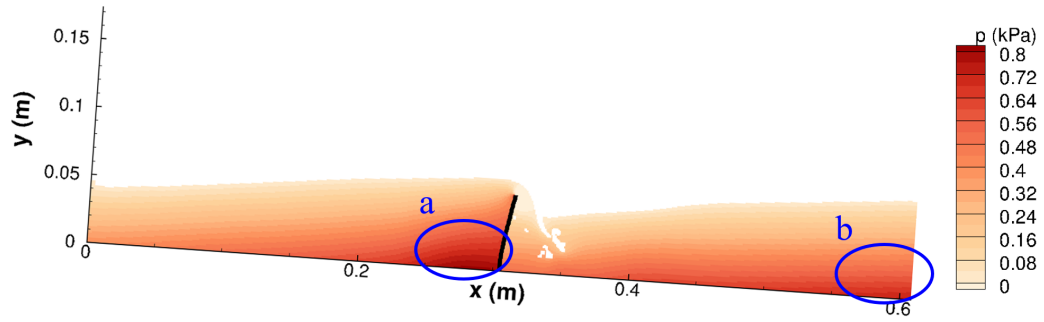
Grid resolution near the plate	Average oscillation amplitude of the plate tip
2.00×10^{-3} m	0.0794 m
1.00×10^{-3} m	0.0854 m
6.67×10^{-4} m	0.0871 m
5.00×10^{-4} m	0.0870 m

1.00×10^{-3} m, and is stretched as the distance to the plate increases. We also tested other Eulerian grid resolutions with $\Delta x = \Delta y = 5.00 \times 10^{-4}$ m, 6.67×10^{-4} m and 2.00×10^{-3} m, and obtained converged result on the average oscillation amplitude of the plate tip for $\Delta x = \Delta y \leq 6.67 \times 10^{-4}$ m, as shown in Table 2.3. The plate is discretized by a uniform Lagrangian mesh, which has 21 grid nodes along the plate length, and 11 grid nodes along the plate width for 3D simulation.

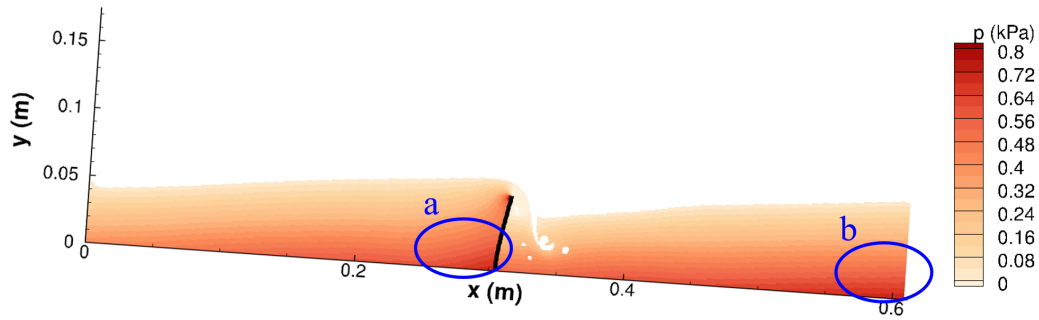
We first study the effect of applying the pressure boundary condition on the plate through the 2D simulation. Figure 2.7 illustrates the pressure fields at $t = 1.7$ s when the pressure boundary condition on the plate is not applied and applied, respectively. As shown in Fig. 2.7(a), when the pressure boundary condition is not applied, the flow pressure on the left side of the plate stem (denoted by “a” in the figure) is higher than that around the bottom-right corner of the tank (denoted by “b” in the figure). However, the pressures at “a” and “b” are expected to be close, because the water depths at these two locations are similar. The difference in the pressures at “a” and “b” is induced by the unphysical pressure gradient near the plate, i.e. in the “a” region. In the case applying the pressure boundary condition on the plate, the flow pressure in the two regions are distributed similarly, as shown in Fig. 2.7(b). Figure 2.8 shows the

profiles of the pressure along the tank bottom in the two cases. According to Eq. (2.8), because the plate stem is the rotation pivot, $D\mathbf{u}_\Gamma/Dt = 0$. On the left surface of the plate, $\mathbf{g} \cdot \mathbf{n} < 0$, therefore $\partial p/\partial n|_\Gamma < 0$, which means that the pressure along the tank bottom should increase when the plate is approached from the left side. It is seen that in the case applying the pressure boundary condition, the pressure jump across the plate is sharp, and the pressure gradient is positive on the left surface of the plate, which agrees with the pressure boundary condition. In the case not applying the pressure boundary condition, the pressure profile is flat adjacent to the left surface of the plate, and $\partial p/\partial n|_\Gamma > 0$ on the left surface of the plate, which violates the pressure boundary condition.

Next, we compare the results of the 3D and 2D simulations. In the 3D simulation, the difference of the horizontal displacement along the free edge never exceeds 0.06 mm, indicating that the deformation of the plate remains 2D. Figure 2.9 shows the time history of the horizontal displacement of the free plate edge for the present 2D and 3D simulations, compared with the experiment and 2D simulation results of [4]. When $t < 1.2$ s, our 2D simulation result agrees with the previous 2D simulation and our 3D simulation result agrees with the experiment, which validate our simulation results. Moreover, the horizontal displacement of the free plate edge in the 3D simulation is smaller than that in the 2D simulation when $t < 1.2$ s and when they are around their peak values at $t = 1.6$ s, 2.5 s, 3.3 s and 4.2 s. This is because in the 3D simulation, the water can go through the gap between the plate and the lateral wall, such that the



(b) Without pressure boundary condition



(a) With pressure boundary condition

Figure 2.7: Comparison of the pressure field at $t = 1.70s$ from the 2D simulation of oil tank sloshing when (a) the pressure boundary condition on the plate is applied and (b) the pressure boundary condition is not applied.

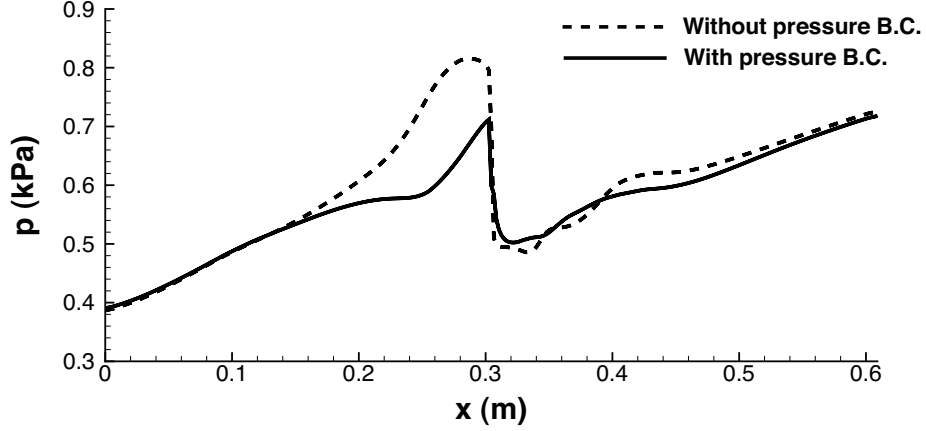


Figure 2.8: Profiles of the pressure along the tank bottom at $t = 1.70$ s for the oil tank sloshing case with and without the pressure boundary condition applied. The plate is located at $x = 0.3045$ m.

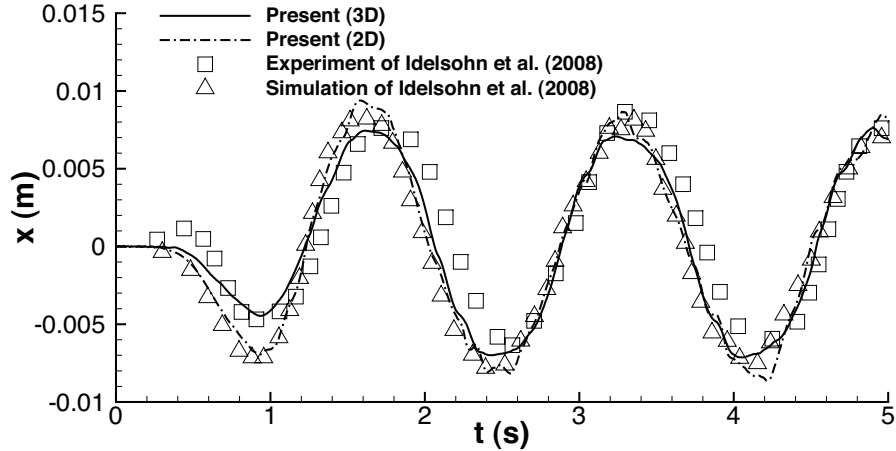


Figure 2.9: Time history of the horizontal displacement of the top edge of the plate in the oil tank sloshing case. Our 2D and 3D simulation results are compared with the experiment and simulation results of [4].

hydrodynamic impact to the plate is weakened compared with the 2D simulation.

Figure 2.10 shows snapshots of the free surface shape and the plate deformation in the middle of the tank span ($z = L_z/2$) and at the front face ($z = L_z$) in the 3D

Table 2.4: Parameters of the case of water impinging onto a plate.

Variables	Values
Density of water ρ_w	$1.0 \times 10^3 \text{ kg/m}^3$
Density of air ρ_a	1.29 kg/m^3
Kinematic viscosity of water ν_w	$1.0 \times 10^{-6} \text{ m}^2/\text{s}$
Kinematic viscosity of air ν_a	$1.5 \times 10^{-5} \text{ m}^2/\text{s}$
Thickness of the plate t	0.012 m
Young's modulus of the plate E	$1.0 \times 10^6 \text{ Pa}$
Density of the plate ρ	$2.5 \times 10^3 \text{ kg/m}^3$

simulation. At the section of $z = L_z/2$, the entrainment of a large bubble induced by the hydraulic jump is observed at $t = 1.68 \text{ s}$ when the tank reaches its maximum rotation angle. This phenomenon agrees with the experiment result of [4]

2.3.3 Collapsing of water column and its impinging onto a flexible bottom-fixed plate

The case simulated by [5] is a 3D modification of the canonical 2D benchmark case about the collapsing of water column impinging onto a flexible plate [36, 58]. [5] considered a plate with finite thickness to avoid the numerical instability induced by the discontinuous of the pressure. We use this case but, without the plate thickness resolved, to examine the capability of our FSI computational framework for handling the pressure discontinuity in the 3D situation. As shown in Fig. 2.11, the water is initially confined in a reservoir, and a flexible plate is located in front of the reservoir with one end mounted at the tank bottom and the other end free. The geometry is listed in Fig. 2.11 and the values of the physical properties of the fluid and the structure are listed in Table 2.4.

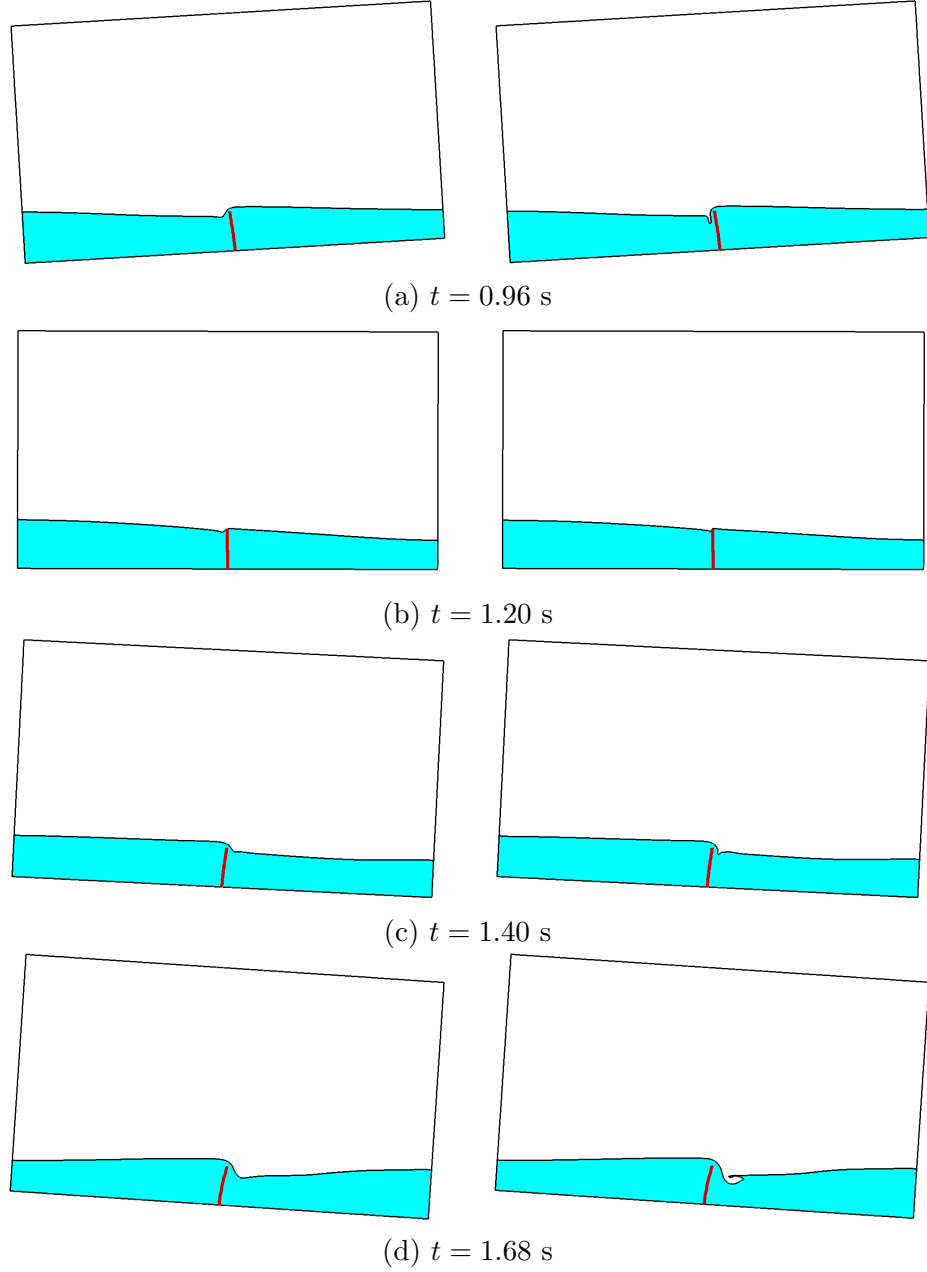


Figure 2.10: Snapshots for the free surface shape and flexible plate deformation in our 3D simulation of the oil tank sloshing case. The upper row shows the section at the front tank wall ($z = L_z = 0.039$ m) and the lower row shows the section at the middle of the tank width ($z = L_z/2 = 0.0195$ m).

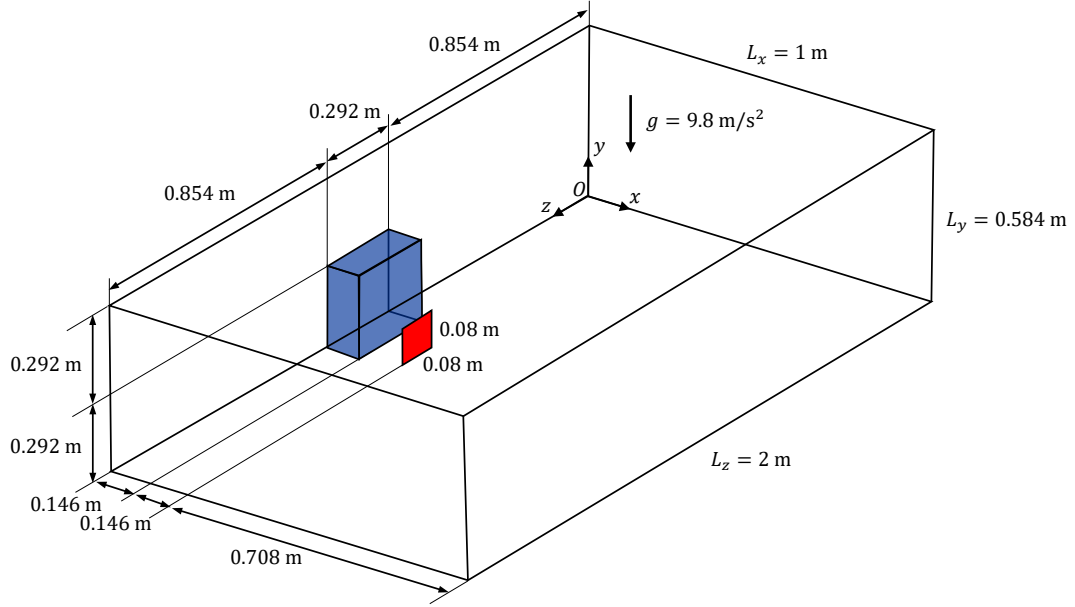


Figure 2.11: Case set-up of collapsing of water column and its impinging onto a flexible bottom-fixed plate.

Note that [5] uses open boundaries on the walls $x = L_x, z = 0$ and $z = L_z$ to avoid the water bouncing back from those walls and influencing the structure deformation, while we extend the computational domain on those directions for the same aim. All the other walls are non-slip. The grid number for the flow field is $200 \times 120 \times 200$, and the grid clusters around the structure for making the grid resolution be uniformly 0.004 m there. The structure is discretized by a uniform 10×10 grid. In this case, we focus on the first impact, which is also the primary impact, of the water column impinging onto the structure, as the deformation and the internal stress of the structure reach maximum during this process.

Figure 2.12 compares the water level during the primary impact, obtained from the

simulations with and without applying the pressure boundary condition. As shown in Fig. 2.12(a-c), water leakage takes place due to the unphysical flux across the plate when the pressure boundary condition is not applied. Note that the leakage is more prominent at the bottom of the plate, because the pressure difference across the plate is higher there than the other places. The water leakage issue is fixed by applying the pressure boundary condition, as shown in Fig. 2.12(d-f). Figure 2.13 compares the time history of the horizontal displacement of the plate tip obtained from the simulations with and without applying the pressure boundary condition and from the simulation by [5]. It is shown that the result from our simulation with the pressure boundary condition agrees with the simulation result by [5]. For our simulation without the pressure boundary condition, the maximum deformation is far smaller than that with the pressure boundary condition, and the different between them is more prominent than that in the 2D case reported in Section 2.3.1. In summary, the above simulation results indicate that applying the pressure boundary condition can mitigate the unphysical flux and improve the result of plate deformation.

2.4 Discussions

The proposed algorithm applies a pressure boundary condition on the immersed boundaries of thickness-unresolvable structures. The boundary conditions are defined on each of the two sides of a thickness-unresolvable structure, and consequently, the flow on one

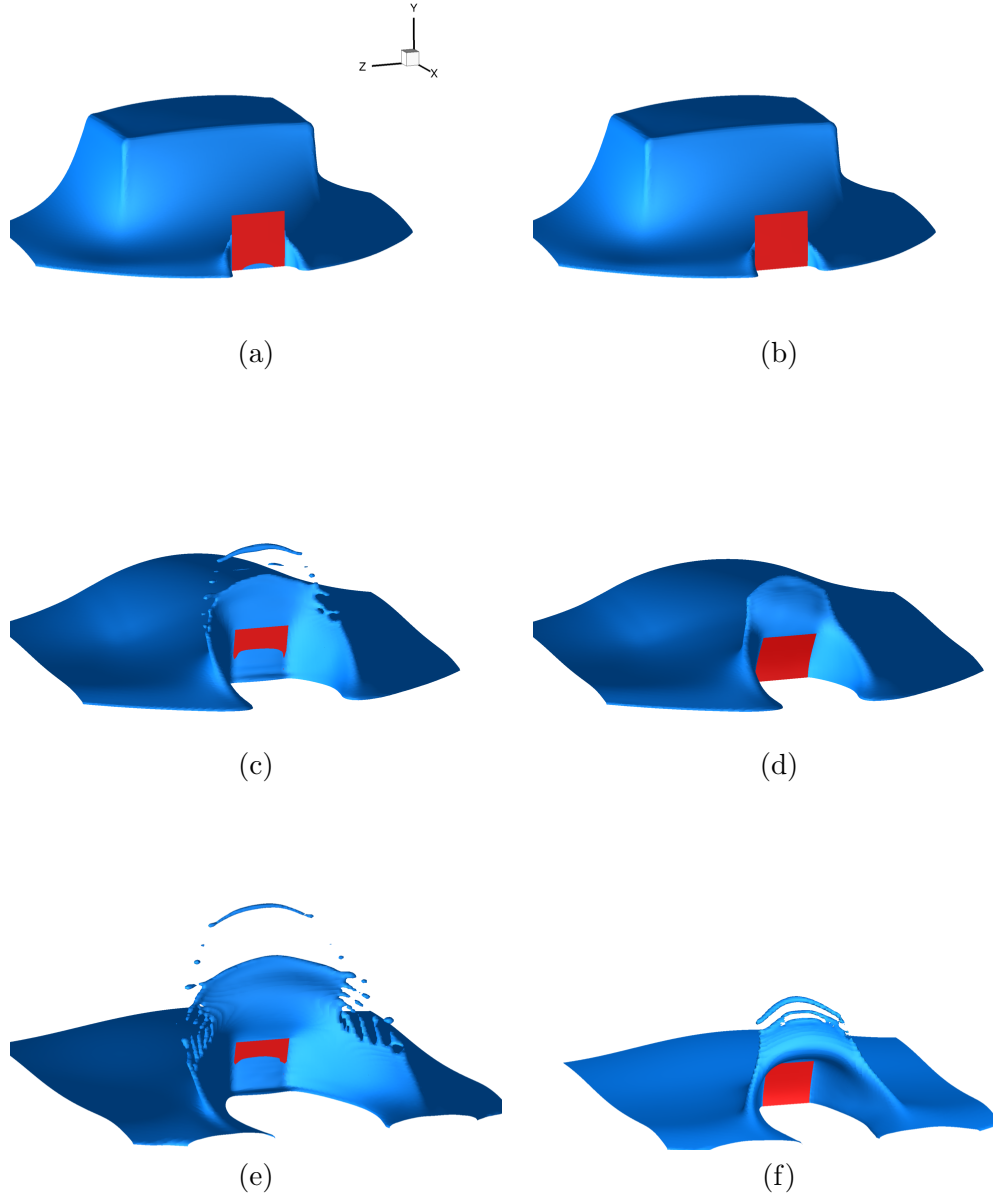


Figure 2.12: Snapshots of the simulation for water impinging onto a plate during the primary impact. (a,c,e) The pressure boundary condition on the immersed boundary is not applied. (b,d,f) The pressure boundary condition on the immersed boundary is applied. (a-b) $t = 0.16$ s, (c,d) $t = 0.24$ s, (e,f) $t = 0.32$ s.

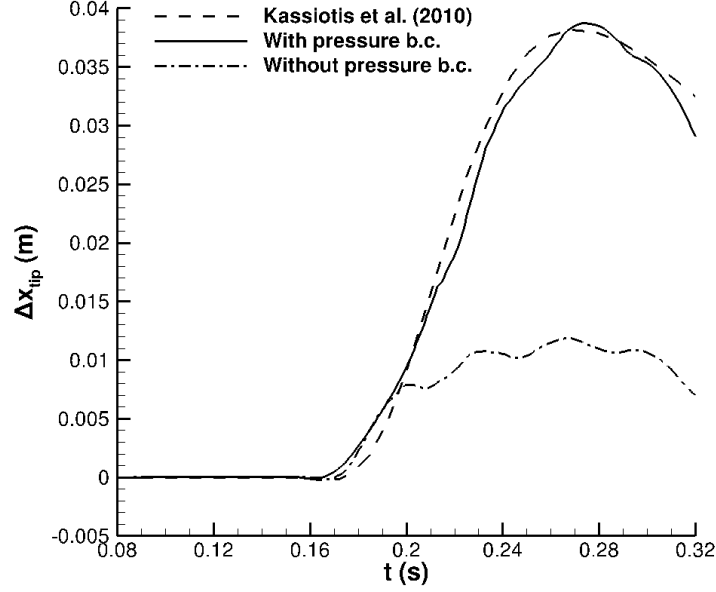


Figure 2.13: Horizontal displacement of the plate tip for the 3D simulation of water impinging onto a plate during the primary impact. The results for applying and not applying the pressure boundary condition on the immersed boundary are compared with the simulation result of [5].

side of the structure does not directly interact with that on the other side. In contrast, in the traditional IB method without explicit pressure boundary condition on the structure surfaces, the pressure gradient across the structure unphysically induces the velocity at the cell-surface nodes adjacent to the structures by Eq. (2.6) and induces the water spurious leakage issue.

In order to show the above point more clearly, we tested both the previous algorithm without the pressure boundary condition and our proposed pressure boundary condition in a modified case of the collapsing water column deforming a top-fixed flexible gate

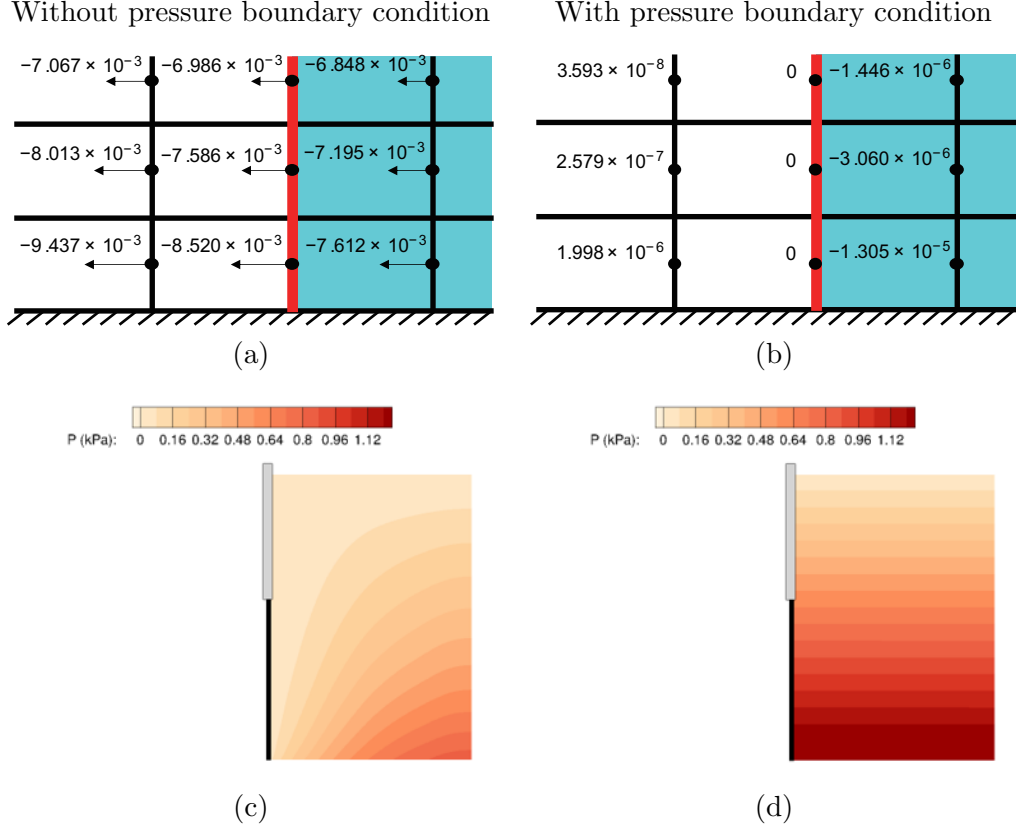


Figure 2.14: Comparison of the results at the first timestep of the modified case of collapsing water column deforming a top-fixed flexible gate (Section 2.3.1), obtained by the algorithms with and without pressure boundary condition on the immersed boundary. The streamwise velocity components on the cell-surface nodes near the gate bottom are shown in (a) and (b), and the pressure fields of the water column are shown in (c) and (d). (a) and (c) are obtained by the algorithm without the pressure boundary, and (b) and (d) are obtained by the algorithm with the pressure boundary.

reported in Section 2.3.1. The case set-up is the same as that shown in Fig. 2.2, except that the rubber gate is fixed to exclude the interference of structure motion. Therefore, the fluid velocity is zero and the water pressure follows the hydrostatic pressure $p_w = \rho g(H - h)$, where h is the water depth. Figure 2.14 shows the results after running the simulation for one timestep, where the streamwise velocity components on the cell-surface nodes near the gate bottom are shown in Figs. 2.14(a) and (b), and the pressure fields of the water column are shown in Figs. 2.14(c) and (d). Figure 2.14(a) shows that when the pressure boundary condition is not applied, the streamwise velocity near the gate are negative, which implies water leakage across the gate. In the case with the pressure boundary condition, shown in Fig. 2.14(b), the streamwise velocity on the plate is exactly zero, and the streamwise velocities at the other locations are three or more orders of magnitude smaller than in Figs. 2.14(a). This small velocity residual further decreases with the increasing iteration steps for solving the Poisson equation, indicating that the non-zero velocity residual is caused by numerical error in the solution of the Poisson equation. By comparing Figs. 2.14(c) and (d), it is seen that the pressure field is similar to that in a dam-break case (for example, Fig. 26(a) in [59]) when the pressure boundary condition is not applied, while the pressure field is hydrostatic pressure when the pressure boundary condition is applied. These phenomena can be explained by analyzing the process of the fractional-step method given by Eqs. (2.3)–(2.6). As the initial velocity is zero everywhere for both the fluid and the gate, the intermediate streamwise velocity $\hat{\mathbf{u}}$ calculated from Eq. (2.3) is also zero for both cases. Note that the

velocity boundary condition is only considered in the velocity interpolation step given by Eq. (2.4), and after that the velocity boundary condition is not used. Therefore, if the pressure boundary condition on the immersed boundary is not considered in Eqs. (2.5) and (2.6), the pressure distribution becomes the same as that in a case without the gate at the end of the first time step, and this is the reason that both Figs. 2.14(a) and (c) show results similar to the dam-break case. Figures 2.14(b) and (d) are obtained by applying the pressure boundary condition and thus the results are physical. Also, Fig. 2.14(c) shows that the pressure gradient along the normal direction of the gate is not zero, which does not comply with the zero pressure gradient predicted by Eq. (2.8), and the pressure gradient is larger when the pressure difference across the gate is higher. Although it is difficult to quantify the exact error of the pressure gradient when the pressure boundary condition is not applied, the present numerical experiment sheds light on the reason causing the violation of the pressure boundary condition in the traditional IB method for thickness-unresolvable structures.

Chapter 3

Canopy Flow Simulation: Introduction, Method, and Case Setup

3.1 Introduction

Research on canopy flows originated from the study of flows interacting with natural organisms such as trees, grasses, and coral reefs ([60, 61, 62, 63, 64, 65]), and its scope has been expanded to many engineering and environmental applications, including wind farms [66] and urban atmospheres [67]. A canopy flow features a mixing layer that develops from the Kelvin–Helmholtz (K–H) instability induced by the difference in velocity

between inside and outside the canopy [68, 69]. In flexible canopies, the coherent turbulent structures in the mixing layer interact with the canopy, thereby propagating the wavy deformation of the canopy, a phenomenon known as *monami* [70], which influences both the transport of momentum into the canopy and the momentum mixing inside the canopy [71, 72].

Nevertheless, the energy flux within the canopy flow remains poorly understood. The presence of the canopy complicates the flow of energy because the spatial inhomogeneity caused by the stems inside the canopy and the correlation between the canopy drag and flow velocity lead to several new terms in the budget equation for the turbulent kinetic energy (TKE) and a dispersive component of the kinetic energy known as dispersive kinetic energy (DKE) [61]. Although a few pioneering studies have attempted to investigate these terms ([73, 60, 74, 61, 75, 76]), the current understanding of the flow field and its underlying mechanisms is still insufficient. For example, owing to the difficulty of experimentally measuring the canopy drag–flow velocity correlation, the role of the waving term, which is associated with this correlation, remains unclear and thus has usually been neglected in previous modeling efforts [73, 61, 60]. In addition, field and laboratory experiments have observed the direct flow of energy from large scales to small scales, also known as the ‘spectral shortcut’ [77, 78, 79, 80], which has been considered to develop turbulence models for canopy flow simulations [81]. However, previous experiments measured only the frequency energy spectra and then derived the wavenumber energy spectra using Taylor’s frozen hypothesis, whereas direct evidence

derived from the spatial energy spectra is lacking. Furthermore, the detailed mechanism responsible for the spectral shortcut remains elusive.

The spectral analysis of the energy budget is a valuable tool for investigating turbulent energy processes, including the production, dissipation, and transport of energy among different locations, across different scales, and among different velocity components [82]. Recently, such spectral analyses have been successfully employed in several channel flow studies. [83] studied the spanwise spectra of energy transfer for a friction Reynolds number $Re_\tau = 1700$. They revealed two scale interaction processes (in addition to the normal energy cascade) in the near-wall region: one is the interaction between the energy-containing scales that generate skin friction, and the other interaction is the downward energy transfer from small to large scales related to the formation of wall-reaching motions containing inactive energy. [84] studied the streamwise and spanwise spectra of energy transfer for $Re_\tau = 5200$ and provided a comprehensive scheme of energy flow: TKE is produced in the streamwise velocity component at streamwise-elongated modes, transferred to the modes with higher spanwise wavenumbers, and redistributed to the vertical and spanwise velocity components by pressure inter-component transport across a broad band of wavenumbers. They also found that the streamwise-elongated modes modulate the near-wall dynamics by wall-normal turbulent transport, which was also discovered by [85]. Both [83] and [84] observed an inverse energy cascade where TKE flows from small to large scales. However, some studies did not detect an inverse energy cascade, such as [86] and [87], who studied

channel flows with $Re_\tau = 210$ and 180, respectively. [88] performed spectral energy budget analyses for turbulent winds over a flat surface, a slow water wave, and a fast water wave and found that waves contribute energy to the dominant wavelength scale close to the wave surface via the production term and enhance the spatial transport of turbulence toward the surface. For the cases with waves, an inverse energy cascade was observed to transport energy from small-scale motions to turbulent scales at the dominant wavelength scale close to the wave surface. In the present study, a spectral analysis of the energy budget of canopy flows is performed for the first time to understand the transfer of turbulent energy in canopies.

In recent years, substantial advancements in canopy flow simulations have been reported. The Reynolds-averaged Navier–Stokes (RANS) equations were first simulated using the k - ϵ , k - ω , and Spalart–Allmaras turbulence models [89, 90, 91, 81], all of which require *a priori* parameter values for the simulation cases. With the increase in computational capabilities in recent years, the large-eddy simulation (LES) technique, which is computationally more expensive but requires less modeling effort, has become quite popular. [92] simulated a free-surface flow over a flexible canopy using a very large-eddy simulation (VLES) where the filter size in the Smagorinsky model was one-tenth of the water depth. More recently, [93] studied the flow and mass transfer in an open-channel canopy flow by applying the dynamic Smagorinsky model and introducing an additional diffusivity to model the effect of stem wakes.

Most previous simulations modeled the canopy as a continuous medium for which

the drag coefficient needed to be prescribed. However, a constant drag coefficient is not applicable to deformable canopies, and the existing drag coefficient prediction models are complex and achieved only limited success for specific cases [94]. Therefore, new simulation approaches that explicitly resolve the flow–canopy interaction are desirable. For instance, the immersed boundary (IB) method solves fluid–structure interactions on fixed Cartesian grids, making it a powerful tool for problems with complex structural geometries and kinematics [12, 95, 11, 8, 96, 9, 97, 98]. Hence, by resolving the stems in the canopy with the IB method, the drag on the stems can be obtained directly; i.e., no *a priori* drag coefficient is required. Moreover, the details of the correlation between the drag force and flow velocity can be obtained to analyse the waving term in the energy budget.

Next, we simulate turbulent aquatic canopy flows with LES and resolve the stems in the canopy with an IB method. We test four cases that cover a broad range of stem flexibilities from rigid stems to oscillatory stems to stems yielding to the flow. Consequently, *monami* is also resolved. Conditional averaging with respect to the stem kinematics is performed to quantify the correlation between the flow features and *monami* wave phase. In addition, the spectral TKE budget is analysed to comprehensively understand the energy fluxes in canopies. In particular, the role of the waving term in the energy transfer in the canopy is resolved, and the canopy drag–flow velocity correlation is studied to understand the mechanism responsible for the waving term. Furthermore, a triadic analysis is conducted to investigate the energy transfer among

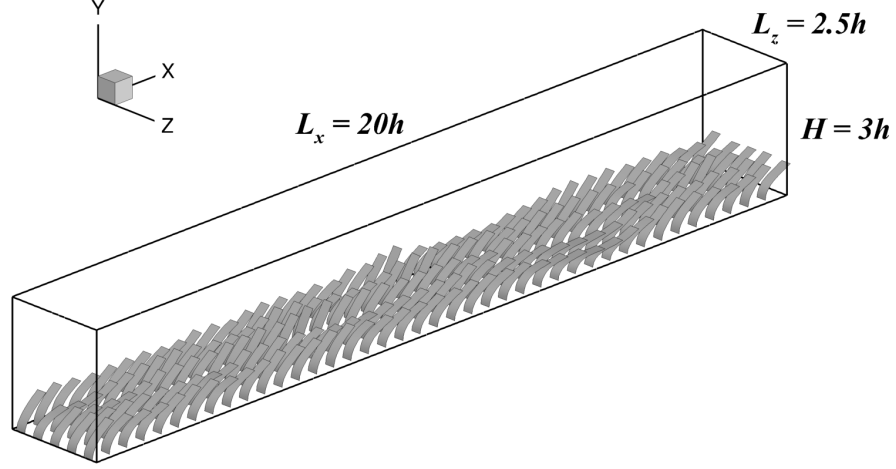


Figure 3.1: Configuration of the simulation setup. The flow is driven by a pressure gradient along the streamwise direction over an evenly spaced 40×5 array of stems. h stands for the height of an undeformed stem.

different scales.

3.2 Method

3.2.1 Flow simulation

In this study, we consider a three-dimensional turbulent flow over a periodic array of flexible canopy stems, as shown in figure 3.1. Details of the dimensions of the computational cases are given in § 3.3. Here, we focus on the mathematical formulation and numerical method. In the Cartesian coordinate system, $x_i (i = 1, 2, 3) = (x, y, z)$, x , y , and z indicate the streamwise, vertical, and spanwise directions, respectively. The resolved velocity components in the LES are denoted by $u_i (i = 1, 2, 3) = (u, v, w)$, which

are governed by the following continuity and momentum equations:

$$\frac{\partial u_i}{\partial x_i} = 0, \quad (3.1a)$$

$$\frac{\partial u_i}{\partial t} + \frac{\partial u_i u_j}{\partial x_j} = \frac{1}{\rho_f} \left(-\frac{\partial p}{\partial x_i} + \mu \frac{\partial S_{ij}}{\partial x_j} + \rho_f g_i - \frac{\partial \tau_{sgs,ij}}{\partial x_j} \right) + f_i, \quad (3.1b)$$

where t denotes time, ρ_f is the fluid density, p is the resolved modified pressure in the LES, μ is the dynamic viscosity of the fluid, $S_{ij} = (\partial u_i / \partial x_j + \partial u_j / \partial x_i) / 2$ is the resolved strain rate tensor, g_i is the gravitational acceleration constant, f_i is the resolved drag force exerted by the canopy, and $\tau_{sgs,ij}$ is the subgrid-scale stress, which is modeled by the dynamic Smagorinsky model [44, 45, 46].

Our flow simulation is performed on a staggered Cartesian grid [99]. Equation (3.1b) is spatially discretized by a second-order central difference scheme and temporally integrated by the second-order Runge–Kutta (RK2) method. During each substep of the RK2 method, the fractional-step approach [100] is employed to enforce the continuity equation (3.1a), where the pressure is solved by a Poisson equation:

$$\frac{\partial}{\partial x_i} \left(\frac{1}{\rho_f} \frac{\partial p}{\partial x_i} \right) = \frac{1}{\alpha \Delta t} \frac{\partial u_i^*}{\partial x_i}, \quad (3.2)$$

where α is 1.0 and 0.5 in the first and second substeps of RK2, respectively; Δt is the timestep of the simulation; and u_i^* is calculated by

$$u_i^* = u_i + \Delta t \left[\frac{1}{\rho_f} \left(\mu \frac{\partial S_{ij}}{\partial x_j} + \rho_f g_i - \frac{\partial \tau_{sgs,ij}}{\partial x_j} \right) + f_i - \frac{\partial u_i u_j}{\partial x_j} \right]. \quad (3.3)$$

Equation (3.2) is solved by the Portable, Extensible Toolkit for Scientific Computation (PETSc) math library [101]. The details of the numerical schemes of our code, its applications to various flows and validations can be found in [40], [102], [41], [103], and [104].

3.2.2 Stem simulation

The stems are modeled as Euler–Bernoulli beams [see, e.g., 48]. The governing equation for the deformation of each stem is

$$\rho_d \frac{\partial^2 X_i}{\partial t^2} = \frac{\partial}{\partial s} \left(T \frac{\partial X_i}{\partial s} \right) - \frac{\partial^2}{\partial s^2} \left(\gamma \frac{\partial^2 X_i}{\partial s^2} \right) + \rho_d g_i + F_i, \quad (3.4)$$

where s is the local coordinate along the stem, $X_i = X_i(s, t)$ is the displacement of the stem in the i direction, T is the tension coefficient, γ is the bending rigidity, F_i is the hydrodynamic force on the stem, and $\rho_d = \rho_s - \rho_f$, where ρ_s is the density of the stem. For a stem whose cross-section is rectangular, $\gamma = Ed^2$, where E is Young’s modulus and d is the stem thickness. In (3.4), the left-hand side (LHS) term is the summation of the stem inertia force and Froude–Krylov force. The first three terms on the right-hand side (RHS) are successively the tension force, the bending force, and the summation of the gravity and buoyancy forces. The Froude–Krylov force and the buoyancy force are modeled in the equation because the stem thickness is too small to be resolved in the

flow simulation. We further assume that the stem is inextensible such that

$$\frac{\partial X_k}{\partial s} \cdot \frac{\partial X_k}{\partial s} = 1. \quad (3.5)$$

To solve the above equations, we first obtain the tension coefficient T from the Poisson equation derived from (3.4) and (3.5):

$$\frac{\partial X_k}{\partial s} \cdot \frac{\partial^2}{\partial s^2} \left(T \frac{\partial X_k}{\partial s} \right) = \frac{1}{2} \frac{\partial^2}{\partial t^2} \left(\frac{\partial X_k}{\partial s} \cdot \frac{\partial X_k}{\partial s} \right) - \frac{\partial^2 X_k}{\partial t \partial s} \cdot \frac{\partial^2 X_k}{\partial t \partial s} - \frac{\partial X_k}{\partial s} \cdot \frac{\partial}{\partial s} (F_{b,k} + F_k), \quad (3.6)$$

where $F_{b,k} = -\frac{\partial^2}{\partial s^2} \left(\gamma \frac{\partial^2 X_k}{\partial s^2} \right)$ is the bending force. Then, the tension coefficient is substituted into (3.4) to solve for X_i . The stem is modeled with a free end and a hinged end. The boundary conditions for the free end are

$$T = 0, \quad \frac{\partial^2 X_i}{\partial s^2} = 0, \quad \frac{\partial^3 X_i}{\partial s^3} = 0, \quad (3.7)$$

and the boundary conditions for the hinged end are

$$X_i = X_{0,i}, \quad \frac{\partial X_i}{\partial s} = \frac{\partial X_{0,i}}{\partial s}, \quad (3.8)$$

where $X_{0,i}(s) = X_i(s, t = 0)$ is the initial condition. Details of the numerical algorithms employed for the stem simulation can be found in [48].

3.2.3 Diffused IB method

The dynamics of the fluid and canopy stems are coupled through the force terms, namely, f_i in (3.1b) and F_i in (3.4), by an IB method [11]. We use the penalization approach developed by [105], where the hydrodynamic force on the structure, F_i , is calculated by

$$F_i(s, t) = \alpha \int_0^t [U_{f,i}(s, \tau) - U_i(s, \tau)] d\tau + \beta [U_{f,i}(s, t) - U_i(s, t)], \quad (3.9)$$

where $U_{f,i}$ is the interpolated fluid velocity on the structure, $U_i = dX_i/dt$ is the Lagrangian velocity of the structure, and α and β are two parameters with dimensions of $ML^{-3}T^{-2}$ and $ML^{-3}T^{-1}$, respectively. Suppose that F_i is the only hydrodynamic force on the structure, $F_i = \rho_f d(U_{f,i} - U_i)/dt$; then, (3.9) describes the behaviour of a mass-spring-damper system about the relative velocity between the flow and the structure, $U_{f,i} - U_i$, as

$$\rho_f \frac{d(U_{f,i} - U_i)}{dt} - \beta (U_{f,i} - U_i) - \alpha \int_0^t (U_{f,i} - U_i) dt' = 0. \quad (3.10)$$

The natural frequency ω_n and the damping ratio of this system ζ are expressed as

$$\omega_n = \sqrt{\frac{-\alpha}{\rho_f}}, \quad \zeta = -\frac{\beta}{2\sqrt{-\alpha\rho_f}}. \quad (3.11a, b)$$

The frequencies of the energetic flow structures must be lower than ω_n to be resolved, which requires a large value of $-\alpha$. However, if the $-\alpha$ value is too large, numerical

instability arises [105]. A moderate α value should therefore be selected to balance the numerical stability and the resolution of the energetic flow structures. Additionally, a large $-\beta$ value can strengthen the numerical stability, but to avoid overdamping, $-\beta$ should not be too large [105]. Details on the selection and testing of the α and β values in our simulations are given in § 3.3.

The interpolated fluid velocity on the stem, $U_{f,i}$, is calculated by integrating the surrounding fluid velocity weighted by a Dirac delta function:

$$U_{f,i}(s, t) = \int_{\Omega} u_i(x_1, x_2, x_3, t) \prod_{j=1}^3 \delta_j (X_j(s, t) - x_j) dx_1 dx_2 dx_3, \quad (3.12)$$

where Ω indicates the computational domain and the Dirac delta function δ_j is defined as

$$\delta_j(r) = \frac{1}{\Delta x_j} \phi \left(\frac{r}{\Delta x_j} \right), \quad \text{no Einstein's summation notation on } j, \quad (3.13a)$$

where the function ϕ is [14]

$$\phi(r) = \begin{cases} \frac{1}{8} \left(3 - 2|r| + \sqrt{1 + 4|r| - 4r^2} \right), & 0 \leq |r| < 1 \\ \frac{1}{8} \left(5 - 2|r| - \sqrt{-7 + 12|r| - 4r^2} \right), & 1 \leq |r| < 2 \\ 0, & |r| \geq 2. \end{cases} \quad (3.14)$$

The delta function is also used to spread the Lagrangian hydrodynamic force on the

structure, $F_i(s, t)$, to the Eulerian force on the background flow, $f_i(x_1, x_2, x_3, t)$:

$$f_i(x_1, x_2, x_3, t) = -\rho_f \int_{\Gamma} F_i(s, t) \prod_{j=1}^3 \delta_j(x_j - X_j(s, t)) ds, \quad (3.15)$$

where Γ indicates all the stems. Note that in this algorithm, the fluid grid needs to be evenly spaced around the stem to conserve momentum [17].

3.2.4 Algorithm validation

We have performed extensive tests on these algorithms. The flow solver has been used in previous studies of various problems ([21, 41, 106, 107]). For the flow–stem interaction problem considered herein, two representative test cases are presented below. The first case is the deformation of a flexible stem in a uniform water current following the experiment of [6], where the stem is hinged on the bottom wall. The Young’s modulus of the stem is $E = 500$ kPa, the stem length h is 0.05 m, the stem width b is 0.01 m, and the stem thickness d is 1.9×10^{-3} m. The current velocity U varies from 0.036 m/s to 0.32 m/s. In our simulation, the grid resolution is $h/40$ for all three dimensions. Figure 3.2 plots the horizontal drag force on the stem as a function of the current velocity. Our result compares well with the measurement data of [6].

The second test case concerns a rigid canopy flow following the experimental study of [7]. A 10×3 array of rigid stems is placed on the bottom, and the computational domain size is $(3.2h, 3h, 1.92h)$ in the streamwise, vertical and spanwise directions. The length h , width b and thickness d of each stem are 0.05 m, 0.008 m and 0.001 m,

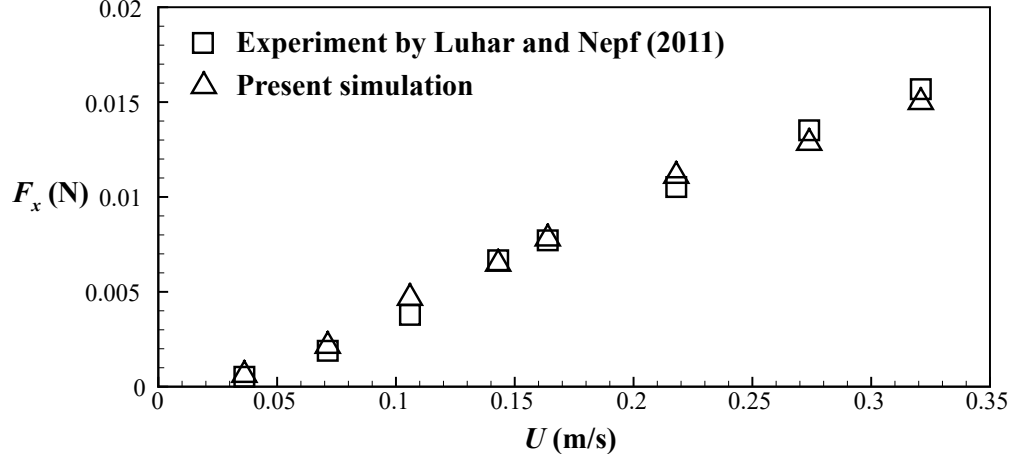


Figure 3.2: Comparison of the horizontal hydrodynamic drag F_x on the stem as a function of the inflow water current speed U in the first validation case of § 3.2.4: $\square$, the result from the experiment by [6]; $\triangle$, the result from the present simulation.

respectively. The flow has a friction velocity at the bottom of $u^* = 0.0253$ m/s. In our simulation, the grid resolution is $h/40$ for all three dimensions. Figure 3.3 compares our vertical profiles of the mean streamwise velocity and the Reynolds shear stress with the measurements of [7]. Both curves agree with the experimental results.

3.3 Case setup

The simulation setup is visualized in figure 3.1. The channel dimensions are $(L_x, H, L_z) = (20h, 3h, 2.5h)$ in the streamwise (x), spanwise (y) and vertical (z) directions, where h is the height of the undeformed stem. An array of aligned stems hinged on the bottom of an open channel is considered. The array has 5 rows parallel to the streamwise direction, and each row has 40 stems. The stem density ρ_s is 1.5 times the fluid density ρ_f , and

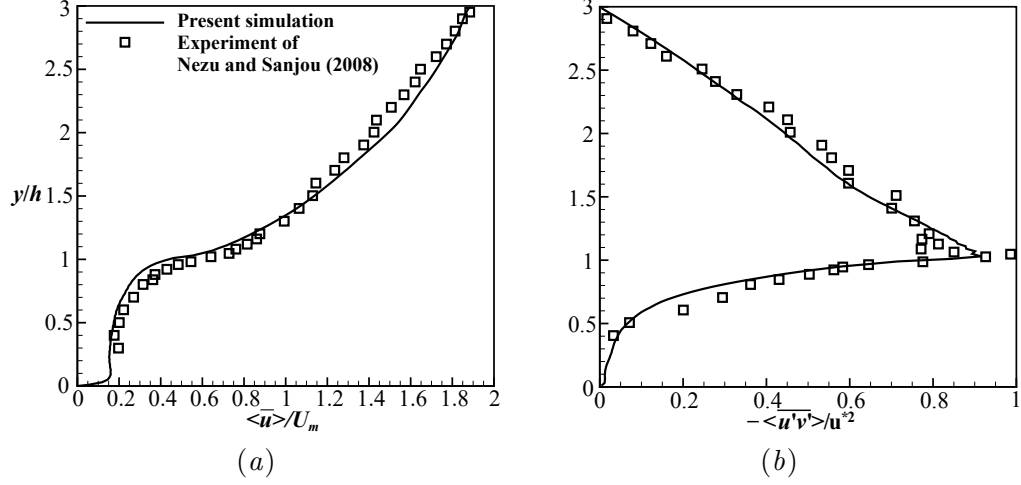


Figure 3.3: Comparison of the horizontally averaged (a) mean streamwise velocity $\langle \bar{u} \rangle$ normalized by the channel mean streamwise flow velocity U_m and (b) Reynolds shear stress $-\langle u'v' \rangle$ normalized by u^{*2} as functions of the vertical height y normalized by the canopy height h in the second validation case of § 3.2.4: $\square$, the result from the experiment by [7]; —, the result from the present simulation.

the stem width b is $0.25h$. The bottom and top boundary conditions of the open channel are nonslip and free-slip, respectively, and the boundary conditions in the streamwise and spanwise directions are periodic. The flow is driven by a constant pressure gradient dp/dx . The friction velocity u^* satisfies $u^* = \sqrt{\frac{1}{\rho_f} \frac{dp}{dx} (H - h)}$. The Reynolds number defined by the canopy height and friction velocity is $Re_\tau = u^*h/\nu = 1000$. The grid resolution Δ is $h/40$ for all three dimensions such that grid number is $800 \times 120 \times 100$. Note that the submergence ratio, the Reynolds number, and the grid resolution in our simulation setup are the same as or close to those in that of [7], whose setup we use to validate our algorithm in § 3.2.4. However, our simulation cases account for stem flexibility, which substantially increases the computational cost.

The stem flexibility is quantified by the dimensionless Cauchy number Ca , which

denotes the ratio between the hydrodynamic force on the stem and the stem restoring force [108]:

$$Ca = \frac{\rho_f u^{*2} b h^3}{EI}, \quad (3.16)$$

where $I = bd^3/12$ is the second moment of area of the stem cross-section. In our simulation cases, we consider $Ca = 0, 5, 30$ and 80 , where $Ca = 0$ corresponds to rigid stems and increasingly larger values of Ca signify stems with greater flexibility. We restrain the stem kinematics to the streamwise-vertical plane and neglect motions such as twisting that cause spanwise displacement. This simplification is based on the following considerations. First, the flow features, such as the coherent vortices in the mixing layer and the classification of different zones in the flow, are dominant in the vertical and streamwise directions, and these features interact with the stem kinematics in these two directions. Second, the spanwise width of the stem is small, so the twisting effect is not prominent. Finally, the streamwise-vertical displacement model for this array of stems is simpler and computationally more affordable than a general three-dimensional displacement model.

For the values of α and β in (3.9), we select $\alpha / [\rho_f (u^*/h)^2] = -10^5$ and $\beta / [\rho_f u^*/h] = -10^2$, respectively, leading to $\omega_n = 316.2u^*/h$ and $\zeta = 0.158$ according to (3.11). This selection ensures that the dominant energetic flow structures are captured and that a reasonably large timestep can be adopted without triggering numerical instability. To compare different values of α and β , we analysed a series of test cases

$\alpha / [\rho_f(u^*/h)^2]$	$\beta / [\rho_f u^*/h]$	d_u/h	d_d/h
-10^3	-10^0	0.802	0.664
$-10^{3.5}$	$-10^{0.5}$	0.787	0.632
-10^4	-10^1	0.731	0.589
$-10^{4.5}$	$-10^{1.5}$	0.730	0.594
-10^5	-10^2	0.737	0.598
$-10^{5.5}$	$-10^{2.5}$	0.727	0.589

Table 3.1: The mean streamwise tip displacements of two tandem stems normalized by the stem length ($d_u/h, d_d/h$) for a range of α and β values. The stems have $Ca = 30$, and their interval follows the simulation setup described in § 3.3.

involving two aligned stems with $Ca = 30$ and the same shape and interval distance as described above. We tested $\alpha / [\rho_f(u^*/h)^2]$ ranging from -10^3 to $-10^{5.5}$ and $\beta / [\rho_f u^*/h]$ ranging from -10^0 to $-10^{2.5}$. Higher absolute values of α and β were not considered because the timestep would be excessively small. The resulting mean streamwise displacements at the tips of the upstream stem, d_u , and the downstream stem, d_d , are listed in table 3.1. The results are insensitive to α and β when their magnitudes are large. Between $(\alpha / [\rho_f(u^*/h)^2] = -10^4, \beta / [\rho_f u^*/h] = -10^1)$ and $(\alpha / [\rho_f(u^*/h)^2] = -10^{5.5}, \beta / [\rho_f u^*/h] = -10^{2.5})$, the relative differences in d_u and d_d are small, both within 2%. Hence, in our simulation cases presented below, $\alpha / [\rho_f(u^*/h)^2] = -10^5$ and $\beta / [\rho_f u^*/h] = -10^2$ are used.

3.4 Definitions of symbols

For the discussion of the results below, we define the following symbols for statistics. For a variable $\phi(x, y, z, t)$, its time average $\bar{\phi}(x, y, z)$ and horizontal spatial average $\langle \phi \rangle(y, t)$

are defined as

$$\bar{\phi}(x, y, z) = \frac{1}{T} \int_0^T \phi(x, y, z, t) dx, \quad (3.17a)$$

$$\langle \phi \rangle(y, t) = \frac{1}{L_x L_z} \int_0^{L_z} \int_0^{L_x} \phi(x, y, z, t) dx dz, \quad (3.17b)$$

where T is the sampling time duration. The time fluctuation $\phi'(x, y, z, t)$ and horizontal spatial fluctuation $\phi''(x, y, z, t)$ are defined as

$$\phi'(x, y, z, t) = \phi(x, y, z, t) - \bar{\phi}(x, y, z), \quad (3.18a)$$

$$\phi''(x, y, z, t) = \phi(x, y, z, t) - \langle \phi \rangle(y, t). \quad (3.18b)$$

By definition, $\bar{\phi}' = 0$ and $\langle \phi'' \rangle = 0$. $\phi(x, y, z, t)$ can be decomposed into a mean component $\langle \bar{\phi} \rangle$, a dispersive component $\bar{\phi}''$ owing to horizontal inhomogeneity, and a turbulent component ϕ' as

$$\phi(x, y, z, t) = \bar{\phi}(x, y, z) + \phi'(x, y, z, t) = \langle \bar{\phi} \rangle(y) + \bar{\phi}''(x, y, z) + \phi'(x, y, z, t). \quad (3.19)$$

Furthermore, $\overline{\langle \phi_1(x, y, z, t) \phi_2(x, y, z, t) \rangle}$ can be decomposed as

$$\overline{\langle \phi_1 \phi_2 \rangle} = \langle \bar{\phi}_1 \rangle \langle \bar{\phi}_2 \rangle + \overline{\langle \bar{\phi}_1'' \bar{\phi}_2'' \rangle} + \overline{\langle \phi_1' \phi_2' \rangle}. \quad (3.20)$$

The RHS of (3.20) consists of the mean term, dispersive term and turbulent term (from left to right). The dispersive term does not exist in a pure channel flow; instead, this term represents the spatial correlation in the time-averaged flow field and exists in the canopy flow owing to the spatial inhomogeneity induced by the stem array [61].

Chapter 4

Overview of Canopy Flow

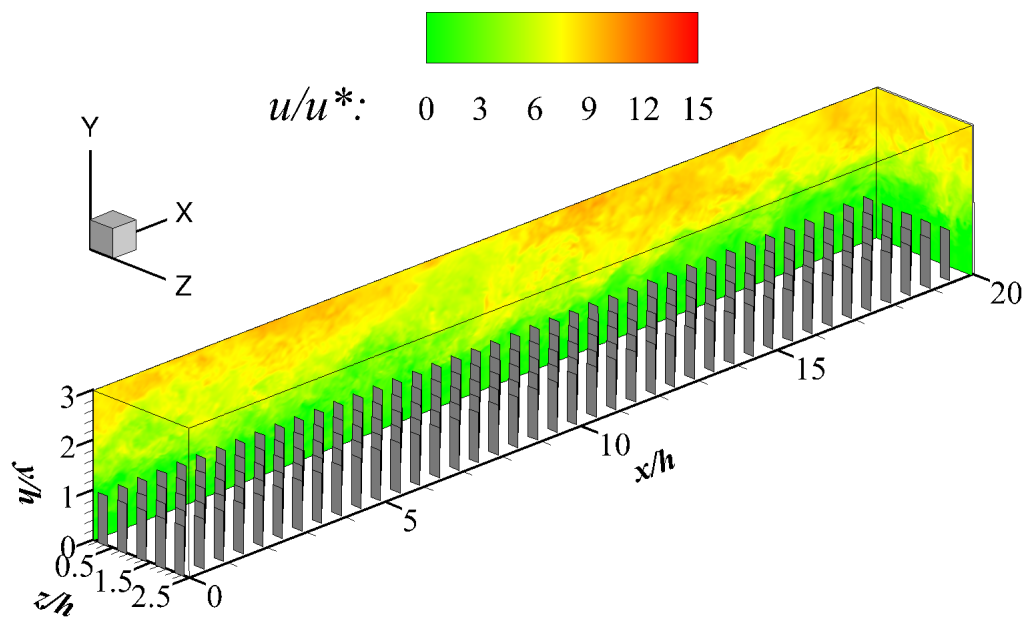
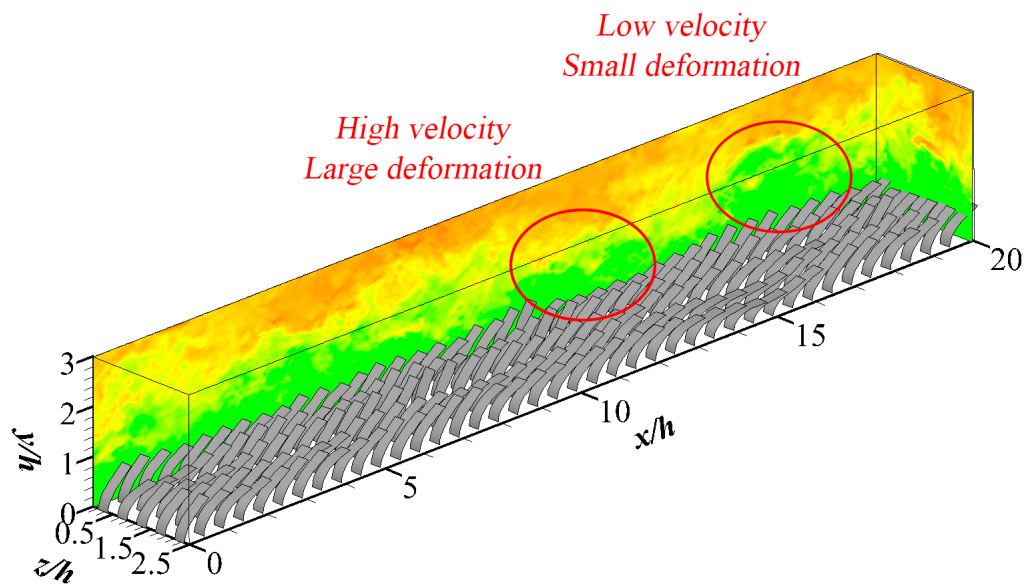
4.1 Instantaneous flow field and canopy deformation

Figure 4.1 shows snapshots of the flow field and canopy deformation for each of the four simulation cases. As the Cauchy number Ca increases, the stem flexibility increases such that the deformation of the canopy stems increases. From the shapes of the deformed stems, we observe that the order of the oscillation mode of the stems increases as the stems become more flexible: for $Ca = 5$, most stems exhibit only the first-order oscillation mode, while for $Ca = 80$, most stems display the second-order oscillation mode. The streamwise velocity in the channel increases as Ca increases because stems with greater flexibility are more prone to reduced drag owing to the reconfiguration of the stem ([109, 110, 111]). For flexible canopies ($Ca \neq 0$), the canopy top has a wave shape that propagates downstream; this phenomenon is known as ‘*monami*’ [70, 71].

The streamwise velocity contours in figure 4.1(*b-d*) reveal coherent structures in the flow. The distribution of these coherent flow structures and the *monami* wave phase appear to correspond to one another. As marked in figure 4.1(*b*), the location where the streamwise velocity is high (low) corresponds to stems with large (small) deformation. These observations are consistent with the common view that *monami* is generated by the interaction between coherent structures in the mixing layer and the deformable canopy [112, 63].

4.2 Velocity and Reynolds shear stress profiles

Figure 4.2 shows the vertical profiles of the mean streamwise velocity $\langle \bar{u} \rangle$ and Reynolds shear stress $-\langle \bar{u}'v' \rangle$. All the mean velocity profiles have inflection points, which is characteristic of the velocity profiles in plane mixing layers [113]. Table 4.1 provides the location of the inflection point for each Ca case. Considering that the inflection point triggers flow instability in the plane mixing layer, we use the location of this inflection point to define the hydrodynamic canopy height, h_c , which we utilize to partition the regions inside and outside of the canopy. In the following discussion of canopy flow, we use the hydrodynamic canopy height as the default canopy height unless otherwise specified. Figure 4.2(*b*) demonstrates that the Reynolds shear stress follows the line $-\langle \bar{u}'v' \rangle = 1 - y/H$ above the maximum value near the hydrodynamic canopy height, indicating that the transport of momentum by molecular viscosity is relatively small compared to the turbulent transport. Below the maximum value, $-\partial \langle \bar{u}'v' \rangle / \partial y > 0$

(a) $Ca = 0$ (b) $Ca = 5$

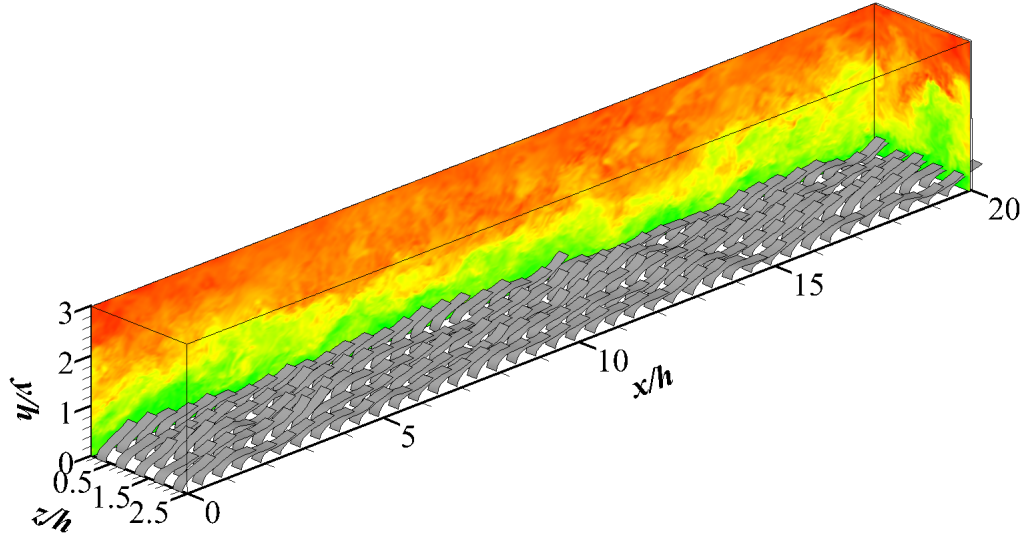
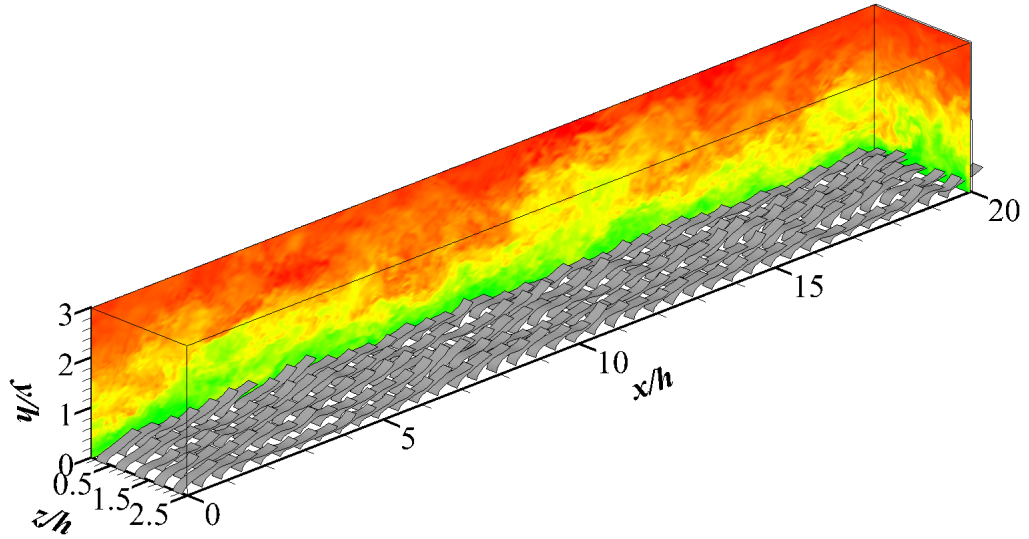
(c) $Ca = 30$ (d) $Ca = 80$

Figure 4.1: Snapshots of the flow field and canopy deformation for $Ca = 0, 5, 30$, and 80 . The x - y and y - z planes show the streamwise velocity u normalized by u^* .

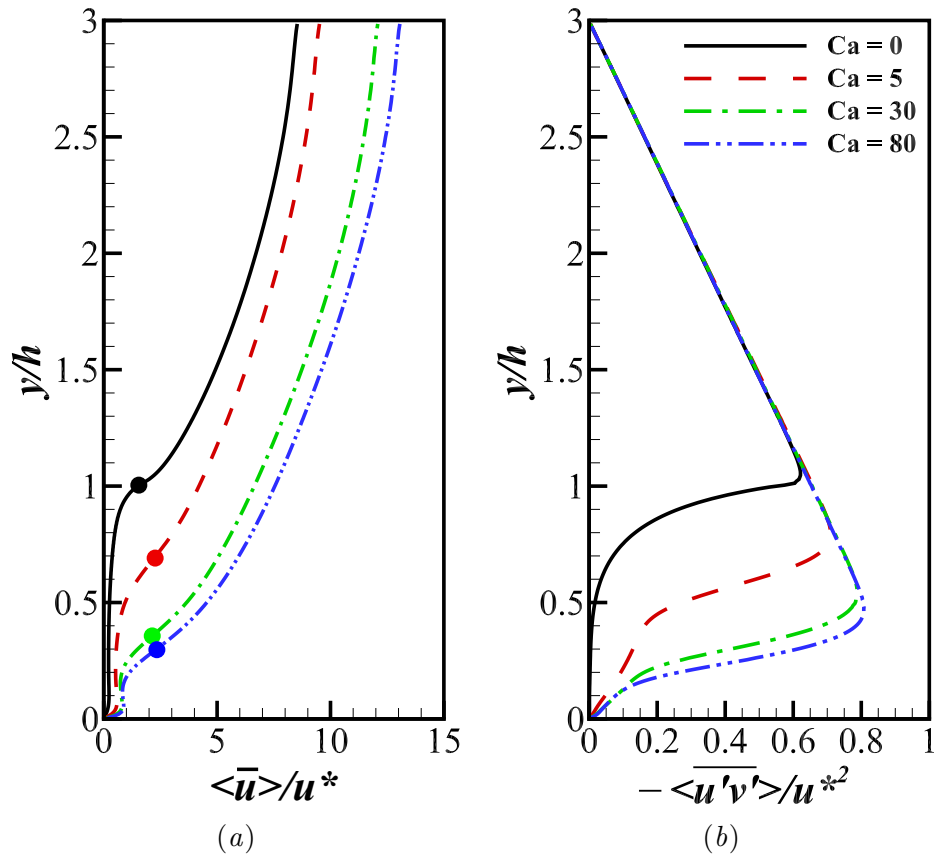


Figure 4.2: Vertical profiles of the (a) mean streamwise velocity $\langle \bar{u} \rangle$ normalized by u^* and (b) Reynolds shear stress $-\langle \bar{u}'v' \rangle$ normalized by u^{*2} for $Ca = 0, 5, 30$, and 80 . The inflection points on the velocity profiles are denoted by dots.

	$Ca = 0$	$Ca = 5$	$Ca = 30$	$Ca = 80$
h_c/h	1.00	0.69	0.36	0.30
h_t/h	1.00	0.76	0.39	0.28
h_{DIS}/h	1.00	0.60	0.31	0.25
h_P/h	1.00	0.71	0.41	0.35
h_{TKE}/h	1.24	0.77	0.55	0.49

Table 4.1: Various characteristic heights in the canopy flow: the height of the inflection point on the mean streamwise velocity profile, h_c ; the average stem tip height, h_t ; the height with the most TKE dissipation, h_{DIS} ; the height with the most TKE production, h_P ; and the height with the most TKE, h_{TKE} . h_c and h_t are also defined as the hydrodynamic canopy height and geometric canopy height, respectively.

owing to canopy drag.

4.3 Stem deformation and *monami*

The average stem tip height h_t is given in table 4.1. Because h_t is convenient to measure experimentally, we use it to define the geometric canopy height. By comparing the stem tip height h_t with the hydrodynamic canopy height h_c in the flexible canopy cases, $h_t - h_c$ decreases as Ca increases. Specifically, $h_c < h_t$ for $Ca = 5$ and 30, and $h_c > h_t$ for $Ca = 80$. This difference can be understood from the fact that compared to the first-order oscillation mode for stems with low flexibility, stems with high flexibility have a second-order oscillation mode where the stem tip is lower.

We also quantify the standard deviations of the stem tip displacements in the horizontal and vertical directions, (σ_x, σ_y) , for flexible canopies. The results are $(\sigma_x/h, \sigma_y/h) = (0.162, 0.132)$ for $Ca = 5$; $(\sigma_x/h, \sigma_y/h) = (0.045, 0.089)$ for $Ca = 30$; and $(\sigma_x/h, \sigma_y/h) = (0.034, 0.072)$ for $Ca = 80$. As Ca increases, the stems clearly fluctuate with a smaller

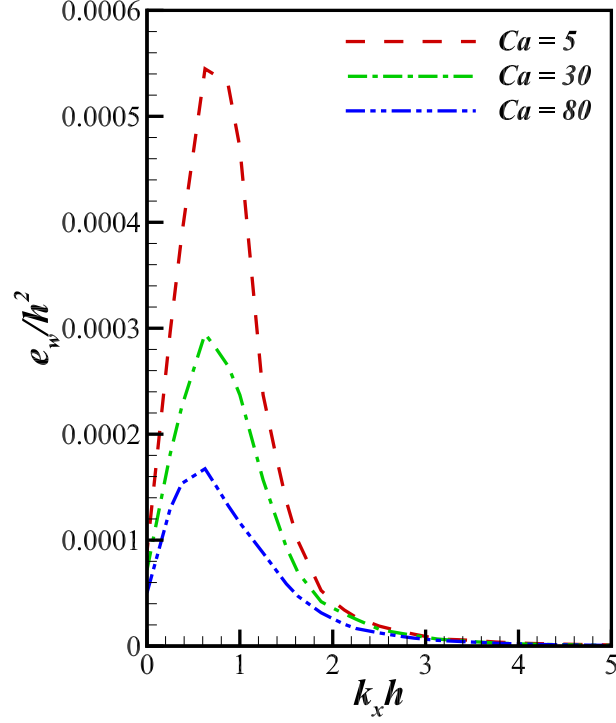


Figure 4.3: Wavenumber spectra of *monami* as defined by (4.1) for cases with $Ca = 5$, 30, and 80.

amplitude, a behaviour that increasingly resembles a rigid stem. This phenomenon occurs because a highly flexible stem yields more to the incoming flow, limiting its deformation fluctuation amplitude. We call this phenomenon ‘high flexibility-induced rigidity’ and use it to explain the similarity between the cases of $Ca = 0$ and 80, which are different from the cases of $Ca = 5$ and 30, in many discussions below.

Next, the wavenumber and frequency–wavenumber spectra of the stem tip envelope are analysed to quantify the kinematic properties of *monami*. We denote the height of the stem tip as a function of the horizontal coordinates and time as $y_t(x, z, t)$; then, its

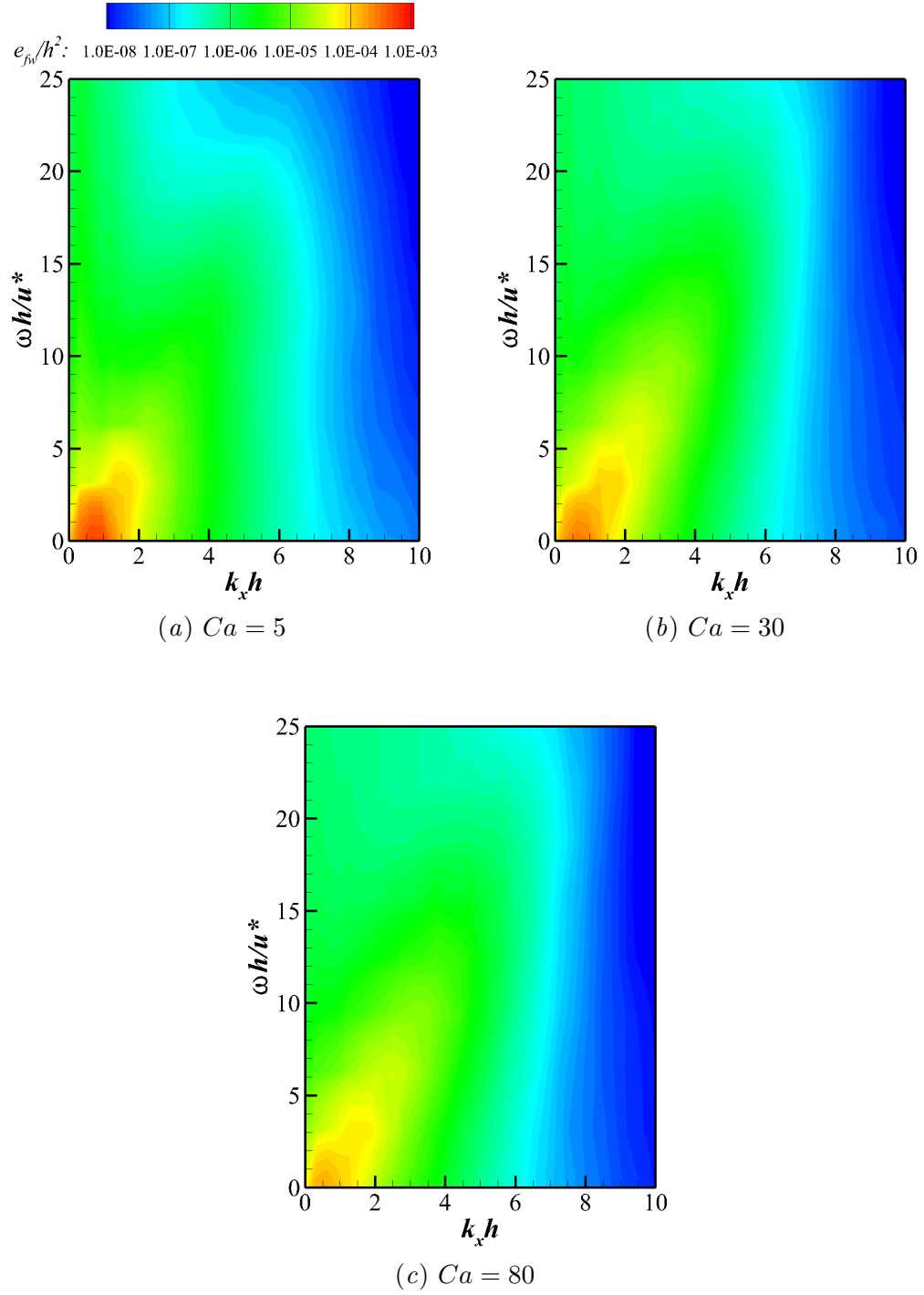


Figure 4.4: Frequency-wavenumber spectra for *monami* as defined by (4.2) for cases with $Ca = 5$, 30 , and 80 .

wavenumber spectrum is

$$e_w(k_x) = \frac{1}{L_x L_z T} \int_0^T \int_0^{L_z} \int_0^{L_x} [y'_t(x, z, t)]^2 e^{-ik_x x} dx dz dt, \quad (4.1)$$

and its frequency–wavenumber spectrum is

$$e_{fw}(k_x, \omega) = \frac{1}{L_x L_z T} \int_0^T \int_0^{L_z} \int_0^{L_x} [y'_t(x, z, t)]^2 e^{-ik_x x} e^{-i\omega t} dx dz dt. \quad (4.2)$$

Figure 4.3 shows the wavenumber spectra. For all cases, the energetic wavenumber components are near $k_x h \sim 0.6$. We define this scale as the ‘*monami* scale’, which corresponds to the scale of the coherent structures in the mixing layer. Likewise, figure 4.4 depicts the frequency–wavenumber spectra for all cases. At every wavenumber, we define a characteristic frequency, $\omega_c(k_x)$, as

$$\omega_c(k_x) = \frac{1}{\int_0^\infty e_{fw}(k_x, \omega) F(k_x, \omega) d\omega} \int_0^\infty \omega e_{fw}(k_x, \omega) F(k_x, \omega) d\omega, \quad (4.3)$$

where

$$F(k_x, \omega) = \begin{cases} 1 & \text{if } e_{fw}(k_x, \omega) > 0.1 \max_{\omega_i} [e_{fw}(k_x, \omega_i)], \\ 0 & \text{otherwise.} \end{cases} \quad (4.4)$$

The extracted *monami* dispersion relation $\omega_c(k_x)$ is plotted in figure 4.5. The relation between the wavenumber and frequency is found to be nearly linear for all cases. The cases of $Ca = 30$ and $Ca = 80$ have similar dispersion relations, i.e., similar wave

phase speeds, because the canopy heights of these two cases are close to each other, while the case of $Ca = 5$ has a lower wave phase speed. Performing linear regression on the dispersion relation curves, we find that the wave phase speeds are $c/u^* = 1.90$ for $Ca = 5$, $c/u^* = 2.55$ for $Ca = 30$, and $c/u^* = 2.57$ for $Ca = 80$. The *monami* phenomenon in all the cases can thus be categorized as a ‘slow wave’ or ‘young wave’ in the context of wind–wave interactions, where the critical layer where $\langle u \rangle - c = 0$ is close to the wave surface and thus does not play a significant role in the flow dynamics; as a result, above the canopy, the flow perturbation by the wave is similar to that by a stationary wavy boundary [see, e.g., 114, 115].

4.4 Conditional averages of the flow velocity and canopy drag

In this section, we illustrate the flow structures associated with the phases of *monami*. Different from water waves with long-lasting wave trains and continuous wave phases, for *monami*, the wave train is short and transient, which makes it difficult to identify the wave phase. Therefore, instead of quantifying the *monami* wave phase and analysing the associated flow field directly, we perform averaging conditioned upon the stem kinematic events corresponding to particular *monami* wave phases. Figure 4.6 illustrates a sketch of the correspondence between the stem kinematic events and the *monami* wave phases. We define four categories of events associated with the stem kinematics: (1) small

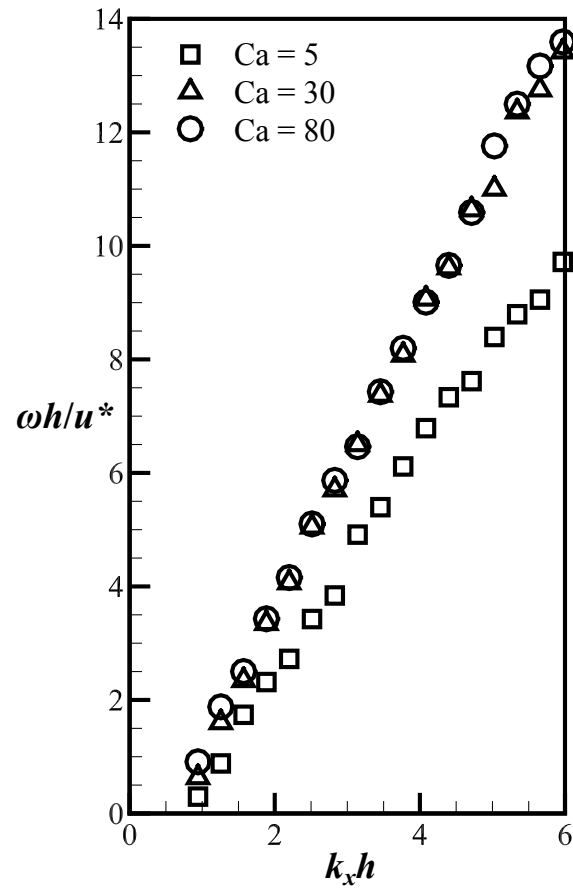


Figure 4.5: Dispersion relation of *monami* for $k_x h \leq 6$ as defined by (4.3) for $Ca = 5$, 30, and 80.

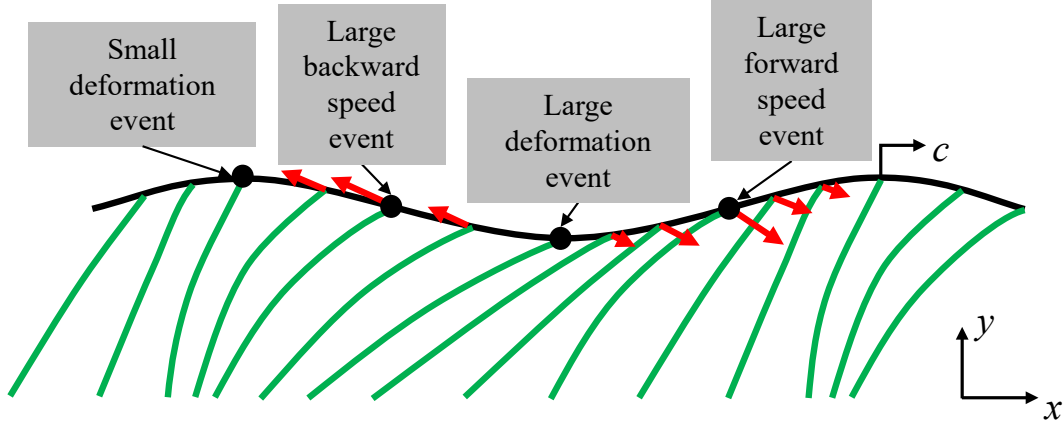


Figure 4.6: Schematics of the relationship between the conditionally averaged events and the *monami* wave phase. The green lines represent the stems. The red arrows represent the velocities of the stem tips. The *monami* is represented by the black envelope connecting the stem tips and travels to the right.

deformation events, which occur when a stem has a large tip height and thus corresponds to the wave crest; (2) large backward speed events, which occur when a stem tip has a large negative horizontal velocity component and thus corresponds to the leeward surface; (3) large deformation events, which occur when a stem has a small tip height and thus corresponds to the wave trough; and (4) large forward speed events, which occur when a stem tip has a large positive horizontal velocity component and thus corresponds to the windward surface. We tested several other definitions, but they yielded similar results; for example, for a large backward speed event, we tested the definitions of a stem tip that has a large positive vertical velocity component and a stem tip that moves backward with a large negative horizontal velocity component.

We perform conditional averages on the above four categories of events and assemble the results to construct an overview of the flow structure associated with the *monami*

phenomenon. The conditional average of a flow variable $\phi(x, y, z, t)$ associated with a stem kinematic event is calculated by

$$\langle \phi \rangle^c(x, y, z) = \frac{1}{N} \sum_{i=1}^N \phi(x_i, y_i, z_i, t_{\text{root},i}), \quad (4.5)$$

where $x_i = (x - x_{\text{root},i}) \bmod L_x$; $y_i = y$; $z_i = (z - z_{\text{root},i}) \bmod L_z$; $(x_{\text{root},i}, z_{\text{root},i})$ and $t_{\text{root},i}$ ($i = 1, \dots, N$) are the root horizontal coordinates and time, respectively, of the stem at which event i occurs; N is the total number of events; and ‘mod’ is the binary modulo operator defined as $a \bmod b = a - b[a/b]$. Note that the modulo operation is performed in the streamwise and spanwise directions because both directions employ the periodic boundary condition, and by performing the modulo operation, the point (x, z) where $x \in (-\infty, 0) \cup [L_x, \infty)$ and $z \in (-\infty, 0) \cup [L_z, \infty)$ can be translated into the computational domain. The above procedure of conditional averaging guarantees that the root of the stem experiencing the concerned event always translates to $(x, z) = (0, 0)$ in the conditionally averaged flow field. The conditional average of the stem geometry $\langle \psi \rangle_{(x,z)}^c(s)$ is calculated as

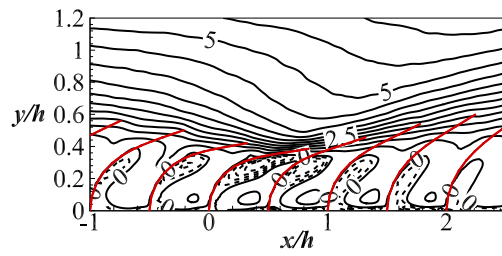
$$\langle \psi \rangle_{(x,z)}^c(s) = \frac{1}{N} \sum_{i=1}^N \psi_{(x_i, z_i)}(s, t_{\text{root},i}), \quad (4.6)$$

where s is the arc-length coordinate and $\psi_{(x,z)}(s, t)$ is the Eulerian coordinate of Lagrangian point s on the stem rooted at (x, z) at time t . To compute the conditional average, we take $N_s = 1500$ time frames of the flow field and select $r_s = 0.5\%$ of all

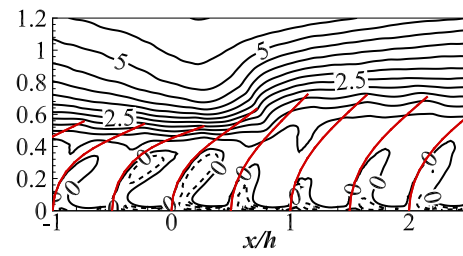
the stem samples for each category of events. Therefore, for each category of events, the total number of samples is $N = N_s r_s N_f = 1500$, where $N_f = 200$ is the number of stems in the flow.

Next, we present the conditional average results. The conditional average of the streamwise velocity component u in the x - y plane that cuts through the centre of the stem row is given in figure 4.7. The relationship between the conditionally averaged events and the *monami* wave phase can be seen from the shape of the conditionally averaged canopy geometry calculated using (4.6). Near the canopy top, the contour lines of u are approximately parallel to the canopy envelope. The trough of the flow streamlines over the stems exhibiting large deformation has a high streamwise velocity, which agrees with the argument that sweep events are associated with canopy deformation during *monami* [116, 72, 117]. Inside the canopy, there are regions with negative u immediately behind stems owing to the backflow, and the size of these regions decreases as Ca increases because of the smaller frontal area. These wake-scale flow structures may impact the transport of scalar quantities and particles, for example, by trapping them to elongate their residence time in the canopy.

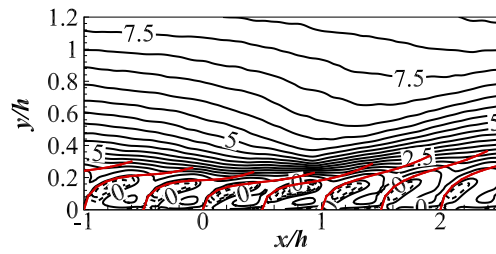
The streamlines of the conditionally averaged flow field in the *monami*-following frame are plotted in figure 4.8 to better illustrate the flow dynamics associated with the *monami* wave phase. Because all the cases with flexible canopies give similar results, only the case of $Ca = 5$ is shown here. The results for the x - y plane cutting through the centre of the stem row are shown in figure 4.8(*a,c,e,g*). From the conditionally averaged

$Ca = 5$


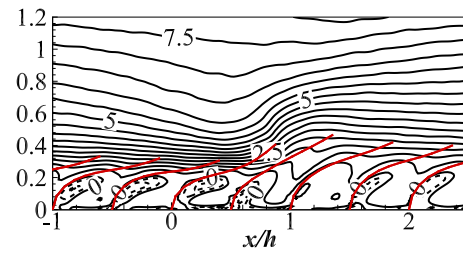
(a)



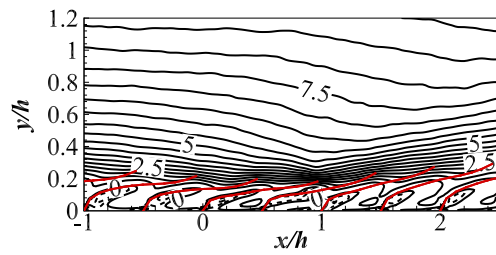
(b)

 $Ca = 30$


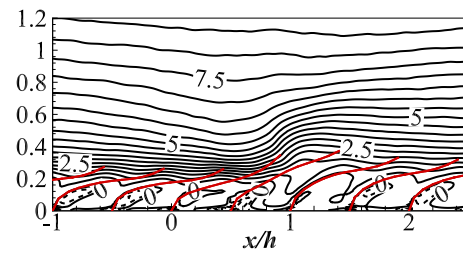
(c)



(d)

 $Ca = 80$


(e)



(f)

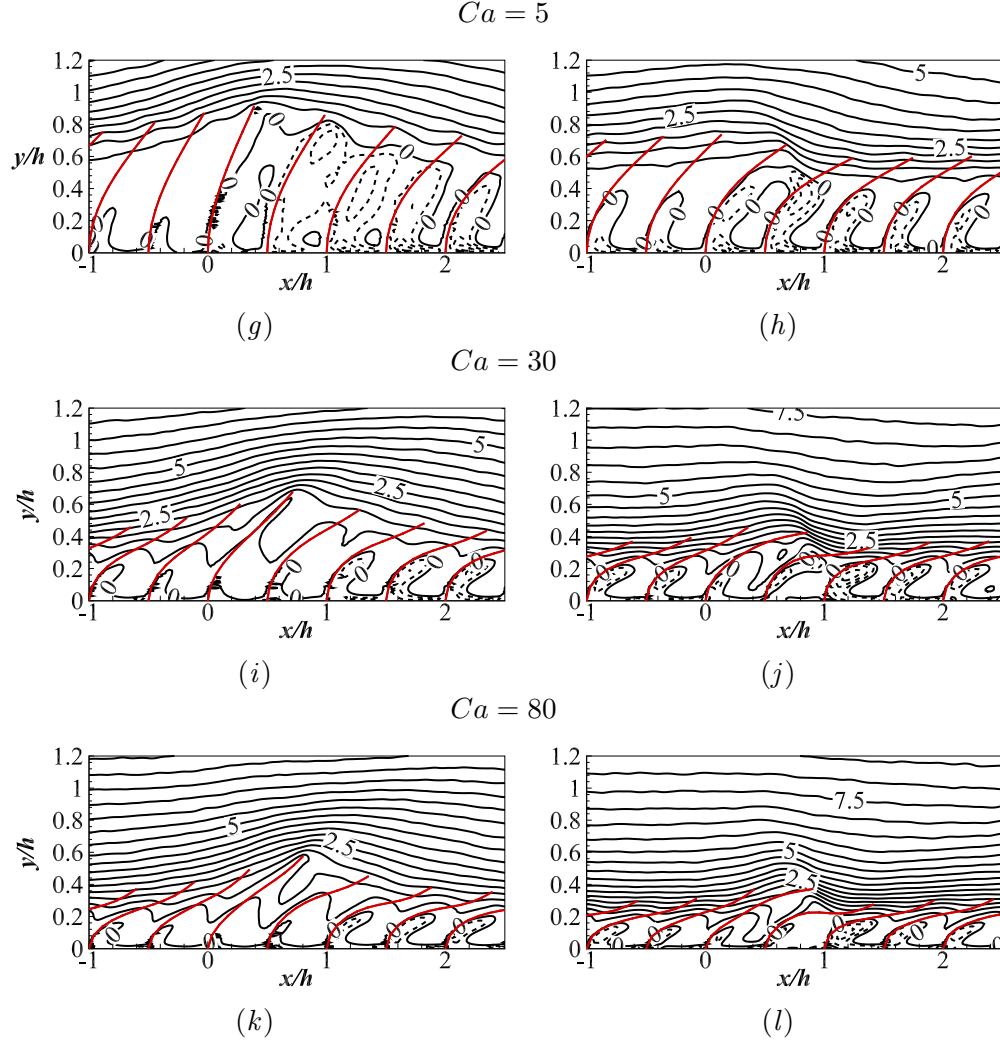


Figure 4.7: Contours of the conditionally averaged normalized streamwise velocity $\langle u \rangle^c / u^*$ in the x - y plane cutting through the centre of the stem row for (a,b,g,h) $Ca = 5$, (c,d,i,j) $Ca = 30$, and (e,f,k,l) $Ca = 80$. The conditionally averaged events are (a,c,e) large deformation events, (b,d,f) large forward speed events, (g,i,k) small deformation events, and (h,j,l) large backward speed events. For the isopleths, the solid and dashed lines represent nonnegative and negative values, respectively, and the intervals between nonnegative values and between negative values are 0.5 and 0.1, respectively. The red lines represent the conditionally averaged shapes of the stems.

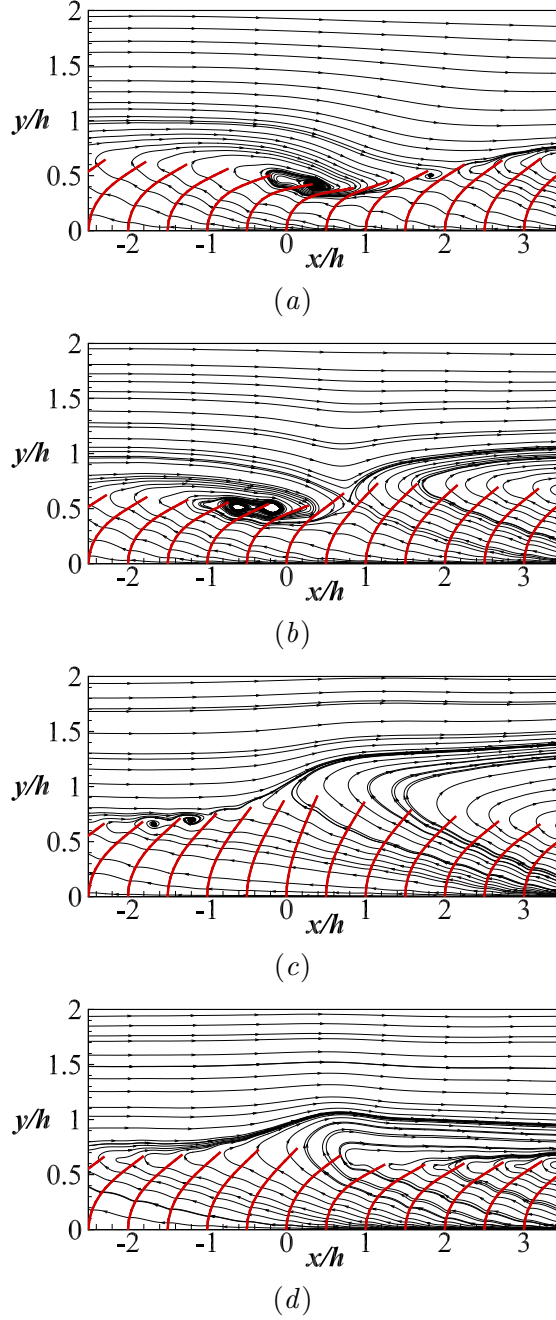


Figure 4.8: Streamlines of the conditionally averaged velocity field in the *monami*-following frame ($(\langle u \rangle^c - c)/u^*, \langle v \rangle^c/u^*$) in the x - y plane cutting through the centre of the stem row. The conditionally averaged events are (a) large deformation events, (b) large forward speed events, (c) small deformation events, and (d) large backward speed events. The red lines represent the conditionally averaged shapes of the stems. The case of $Ca = 5$ is shown.

flow fields for the large deformation and large forward speed events, we observe a vortex residing on the *monami* trough near $x/h = 0$ and $x/h = -0.5$, respectively. This vortex is similar to the cat’s eye structure in wind–wave interactions ([118, 119, 115]), except that the streamlines extend into the canopy owing to the permeability of the canopy. The vortex and the critical layer (where $\langle u \rangle^c - c = 0$) reside near the canopy top, which is consistent with the result in § 4.3 that *monami* is a slow wave. The centre of the vortex, which is a convergence point, resides on the leeward surface. Another large vortex can be observed on the right; note that this vortex is stretched in the streamwise direction, and hence, only part of it is shown in the figure. A saddle point between these two vortices is located on the windward surface, as can be seen in figure 4.8(c), and the boundary between these two vortices can be identified from the saddle point and the influx streamlines connected to it. The streamlines above the vortices deform with them and thus are depressed above the saddle point. The streamlines of the conditionally averaged flow field in the *monami*-following frame are summarized in figure 4.9(a).

Figure 4.10 shows the conditional average of the fluctuations of the streamwise and vertical velocity components, $\langle u' \rangle^c$ and $\langle v' \rangle^c$, respectively. As all the flexible canopy cases show similar results, only the case of $Ca = 5$ is presented. For the streamwise velocity fluctuation u' , we note that the large deformation and large forward speed events are characterized by $u' > 0$. This phenomenon can be understood from two perspectives of the flow–stem interaction. First, from the perspective of the stems affecting the flow, highly deformed stems have smaller frontal areas, and stems with a

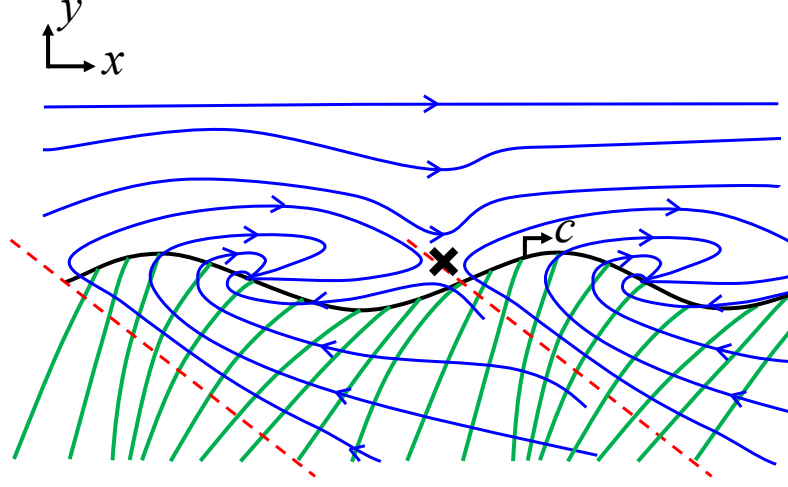


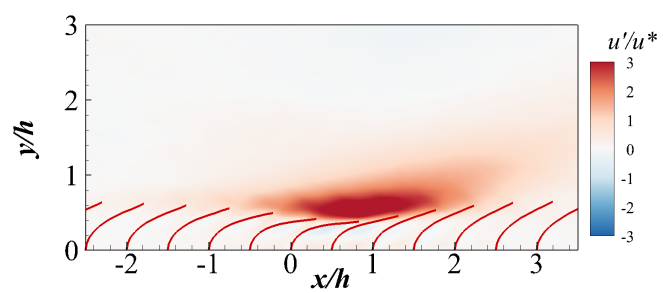
Figure 4.9: Schematics of the flow pattern associated with the *monami* wave phase in the *monami*-following frame in the x - y plane cutting through the centre of the stem row. The blue curves with arrows are streamlines. The red dashed lines indicate the boundaries between the main flow structures. The green lines and black curve represent the stems and the *monami* surface, respectively. The black cross indicates the saddle point.

large forward speed have a smaller velocity relative to the flow; in both situations, the flow resistance decreases, and the flow velocity is increased. Second, from the perspective of the flow affecting the stems, a faster flow can more easily deform the stems and maintain their large deformation. A similar explanation can be applied to $u' < 0$ in small deformation and large backward speed events, where the frontal areas are large and the stems have a large velocity relative to the flow, thereby increasing the drag force, and a slow flow allows the stems to rebound. For the vertical velocity fluctuation v' , the large deformation and large forward speed events are characterized by $v' < 0$, while the small deformation and large backward speed events are characterized by $v' > 0$. In summary, large deformation and large forward speed events are characterized by sweep

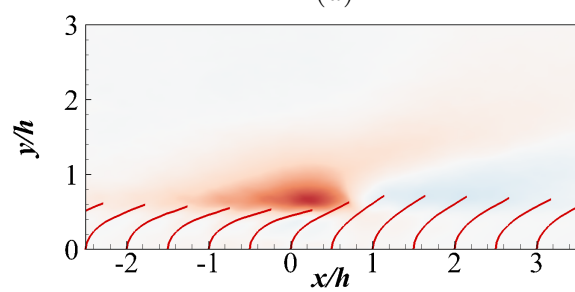
events ($u' > 0, v' < 0$), and small deformation and large backward speed events are characterized by ejection events ($u' < 0, v' > 0$), which is consistent with the findings of previous experiment and numerical studies [116, 72, 117].

Figure 4.11 shows the conditionally averaged fluctuations of the streamwise and vertical components of the drag force, $\langle f'_1 \rangle^c$ and $\langle f'_2 \rangle^c$, respectively. Similar to figure 4.10, only the case of $Ca = 5$ is plotted, as it is representative of the flexible canopies. Within the stem oscillation trajectories, the time average of the streamwise force component $\bar{f}_1$ is always negative because it averages the zero value of f_1 if a stem is not present and the negative value of f_1 if a stem is present. As a result, for the region within the stem oscillation trajectories, f'_1 is positive if a stem is not present because $f_1 = \bar{f}_1 + f'_1$, $f_1 = 0$ and $\bar{f}_1 < 0$. When a stem is present, f'_1 is usually negative. For example, for small deformation events (figure 4.11e), f'_1 is negative around the stem tips in the centre of the domain. This is because this location is the highest reachable location of a stem, so $f_1 < 0$ only for small deformation events and $f_1 > 0$ otherwise. In contrast, for large deformation events (figure 4.11a), f'_1 is negative at the upper part of the stems because the streamwise velocity there is higher than in other events (see figure 4.10), and therefore, the drag force has large negative values.

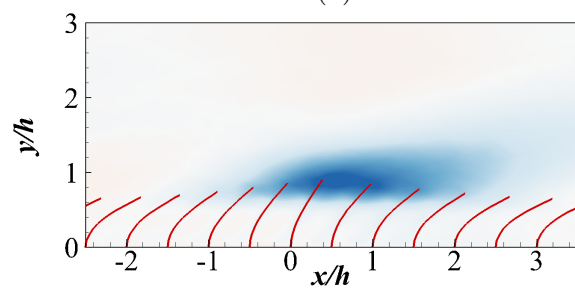
For the vertical force component, its time average $\bar{f}_2$ is positive within the stem oscillation trajectories because the inclination of the stems turns the streamwise incoming flow to the direction along the stem, with a positive vertical velocity component. Then, f'_2 is negative if a stem is not present because $f_2 = \bar{f}_2 + f'_2$, $f_2 = 0$ and $\bar{f}_2 > 0$. The



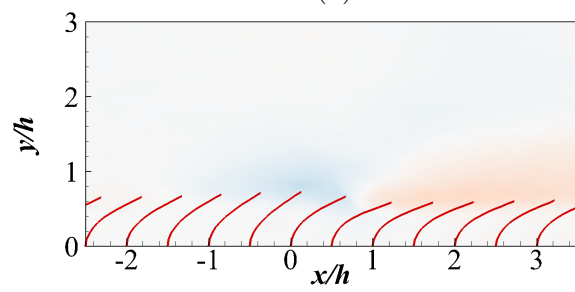
(a)



(b)



(c)



(d)

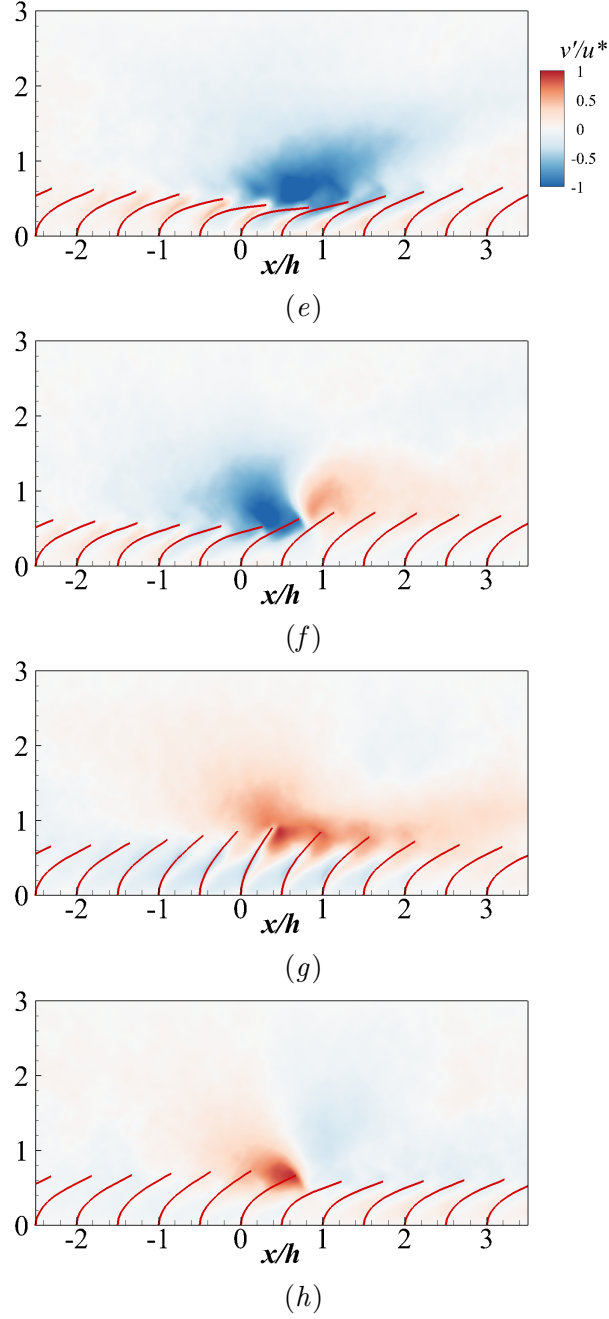
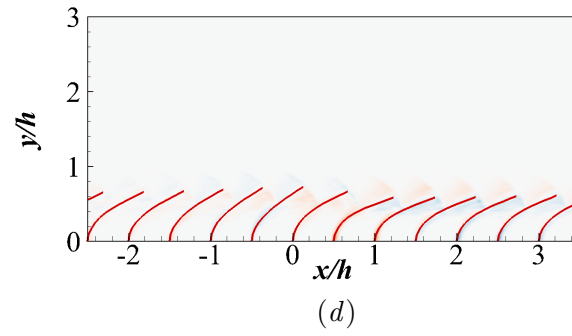
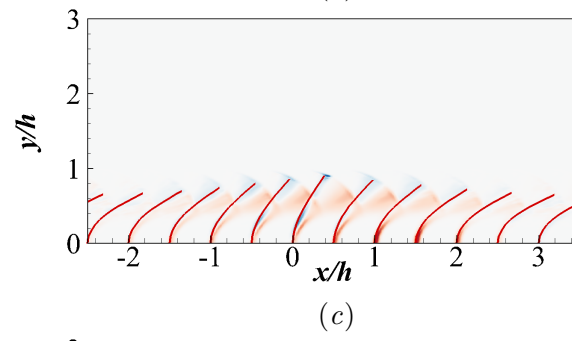
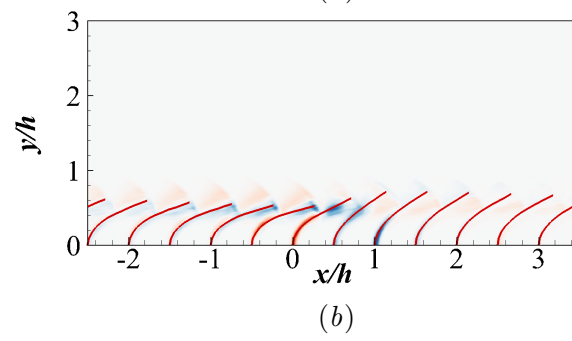
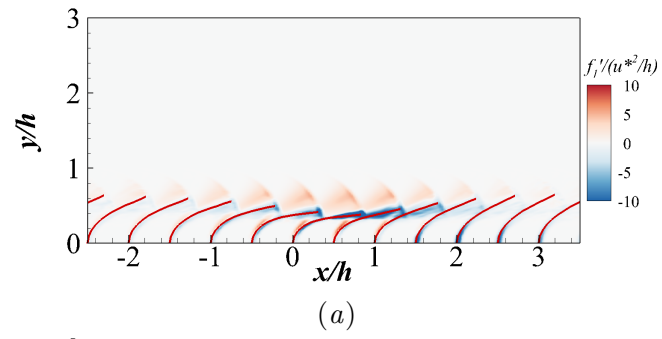


Figure 4.10: Contours of the conditionally averaged ($a-d$) normalized streamwise velocity fluctuation u'/u^* and ($e-h$) normalized vertical velocity fluctuation v'/u^* in the x - y plane cutting through the centre of the stem row. The conditionally averaged events are (a,e) large deformation events, (b,f) large forward speed events, (c,g) small deformation events, and (d,h) large backward speed events. The red lines represent the conditionally averaged shapes of the stems. The case of $Ca = 5$ is shown.

sign of f'_2 when a stem is present depends on the vertical velocity v there. For large deformation and large forward speed events, sweep events near the stems have $v' < 0$, which results in $f'_2 > 0$. In contrast, small deformation and large backward speed events have $f'_2 < 0$ near the stem tips and around the upper stem because of the ejection events with $v' > 0$. The results shown in figures 4.10 and 4.11 are important for the analyses of the waving term in the TKE budget in § 5.6.



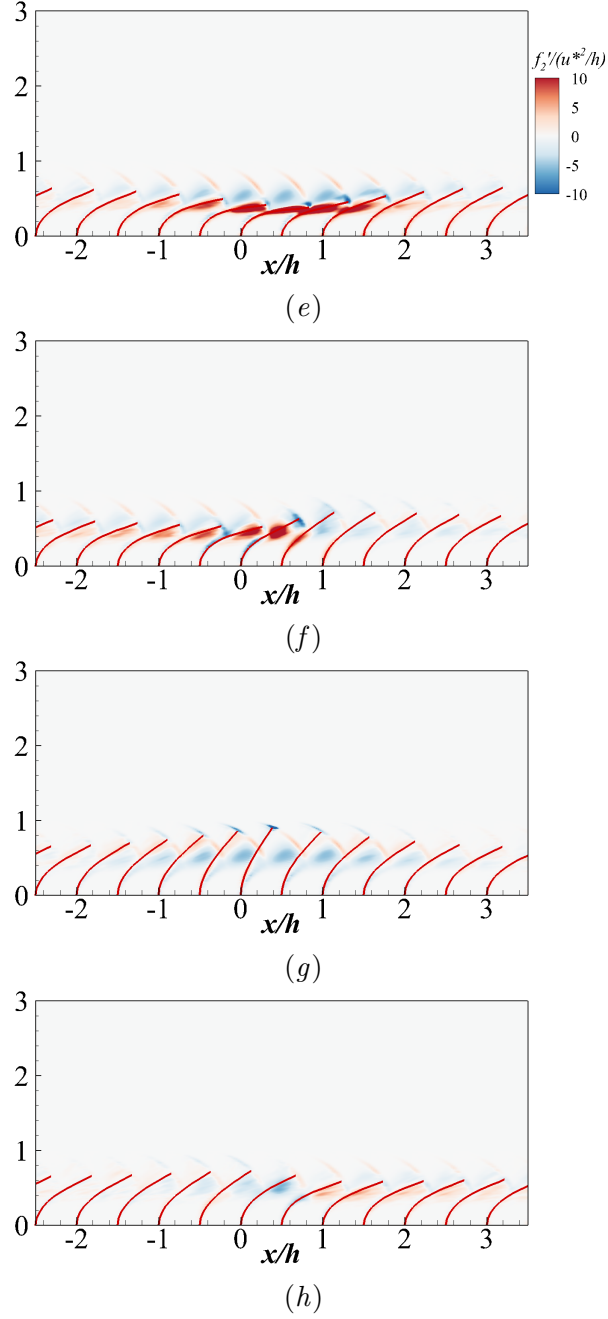


Figure 4.11: Contours of the conditionally averaged ($a-d$) normalized streamwise drag force fluctuation $f'_1/(u^{*2}/h)$ and ($b-h$) normalized vertical drag force fluctuation $f'_2/(u^{*2}/h)$ in the $x-y$ plane cutting through the centre of the stem row. The conditionally averaged events are (a,e) large deformation events, (b,f) large forward speed events, (c,g) small deformation events, and (d,h) large backward speed events. The red lines represent the conditionally averaged shapes of the stems. The case of $Ca = 5$ is shown.

Chapter 5

Spectral Analyses on TKE and DKE Budgets

5.1 Energy transport in the canopy flow: MKE, DKE, and TKE

Following (3.20), the kinetic energy in the canopy flow $e = \langle \overline{u_i u_i} \rangle / 2$ can be decomposed into

$$e = \underbrace{\frac{1}{2} \langle \overline{u_i} \rangle \langle \overline{u_i} \rangle}_{e_{MKE}} + \underbrace{\frac{1}{2} \langle \overline{u_i'' u_i''} \rangle}_{e_{DKE}} + \underbrace{\frac{1}{2} \langle \overline{u_i' u_i'} \rangle}_{e_{TKE}}, \quad (5.1)$$

where e_{MKE} , e_{DKE} , and e_{TKE} are the mean kinetic energy (MKE), DKE, and TKE, respectively. DKE represents the part of the kinetic energy induced by the flow inhomogeneity and does not exist in a pure channel flow. The TKE budget in the canopy

flow is [61]

$$\begin{aligned}
\frac{\partial}{\partial t} \left(\frac{1}{2} \overline{u'_i u'_i} \right) = & \underbrace{-\overline{u'_i u'_j} \frac{\partial \langle \bar{u}_i \rangle}{\partial x_j}}_{P_s} - \underbrace{\left\langle \overline{u'_i u'_j}'' \frac{\partial \bar{u}_i}{\partial x_j} \right\rangle}_{P_w} - \underbrace{\frac{1}{2} \frac{\partial \langle \bar{u}_j'' \overline{u'_i u'_i}'' \rangle}{\partial x_j}}_{T_d} - \underbrace{\frac{1}{2} \frac{\partial \langle \overline{u'_i u'_i u'_j} \rangle}{\partial x_j}}_{T_t} \\
& - \underbrace{\frac{\partial \langle \bar{p}' u'_j \rangle}{\partial x_j}}_{T_p} + \underbrace{\frac{\nu}{2} \frac{\partial^2 \langle \overline{u'_i u'_i} \rangle}{\partial x_j \partial x_j}}_{T_v} - \underbrace{\nu \left\langle \frac{\partial \overline{u'_i}}{\partial x_j} \frac{\partial \overline{u'_i}}{\partial x_j} \right\rangle}_{\epsilon} + \underbrace{\langle \bar{f}' u' \rangle}_W,
\end{aligned} \tag{5.2}$$

where the RHS terms are the shear production term P_s , wake production term P_w , dispersive transport term T_d , turbulent transport term T_t , pressure transport term T_p , viscous transport term T_v , dissipation term ϵ , and waving term W . In these terms, the wake production term P_w and the dispersive transport term T_d are related to the spatial inhomogeneity of the canopy flow, while the waving term W represents the correlation between the hydrodynamic drag and the stem waving motion. In the literature, fewer studies have been conducted on the waving term than on the other terms. Our numerical simulation provides comprehensive descriptions of the hydrodynamic drag and flow velocity in space and time such that the canopy waving effect can be quantified and analysed.

Figure 5.1 plots the vertical profiles of TKE and DKE. Below $y \sim 1.8h$, TKE increases with Ca , and the height of its maximum value decreases as Ca increases because the canopy height decreases. The maximum values of the TKE profiles are above the canopy where the centres of the coherent structures in the mixing layer are located [63]. Above $y \sim 1.8h$, the TKE profiles of all the cases collapse into one

curve, indicating that the TKE enhancement effect of flexible canopies is limited with height. DKE is prominent only in the canopy region. The amplitude of DKE is small compared with that of TKE, which is consistent with the measurement result of [120]. Interestingly, the maximum value of DKE in the rigid canopy is larger than that at $Ca = 5$ and closer to that in the cases of $Ca = 30$ and 80. We believe the reason is that the oscillation amplitude of the stems has a negative correlation with DKE and that highly flexible stems oscillate less and behave more like a rigid stem, i.e. the phenomenon of ‘high flexibility-induced rigidity’ as discussed in § 4.3.

The vertical profiles of the TKE budget terms (5.2) are plotted in figure 5.2. For all cases, the shear production term P_s is the dominant energy source. However, the wake production term P_w can be higher than P_s inside the canopy because of the strong flow inhomogeneity there, which is consistent with the findings of [120] and [75]. The turbulent transport term T_t has the largest magnitude among all four transport terms and plays an important role in balancing the high energy production immediately above the canopy and the high dissipation inside the canopy. In the rigid canopy case, the extrema of P_s and ϵ are $12u_*^3/h$ and $-5u_*^3/h$, respectively, which agree with the particle image velocimetry (PIV) results of [76], who measured a rigid canopy flow in a wind tunnel. The flexible canopy cases have a similar profile shape for each term, except that the terms have larger amplitudes and lower peak heights as Ca increases owing to the lower canopy height and faster flow. Comparing the rigid and flexible canopy cases reveals that the peak of P_s is sharper in the former because the oscillation of stems in

the flexible canopies smears the canopy top and the flow shear there. Another major difference between the rigid and flexible canopy cases is the waving term W . For the rigid canopy case ($W = \langle \overline{f'u'} \rangle = 0$), because $u = 0$ on the stem surface and $f = 0$ away from the stem surface, $u' \neq 0$ and $f' \neq 0$ cannot be satisfied simultaneously. For the flexible canopy cases, the waving term is an energy source near the canopy top where the stem tip has a large oscillation amplitude and a weak energy sink inside the canopy where the stem oscillation is reduced. The maximum value of W is approximately half of P_s , which makes it another major TKE source term for the flexible canopy cases. Owing to the difficulty of experimentally measuring the drag-velocity correlation in flexible canopies, previous studies did not clarify the role of the waving term in canopy flow dynamics [73, 61, 60]. Our simulations, which explicitly resolve the stems, show that the waving term can play an important role in TKE generation in flexible canopies. More analyses of the waving effects are provided in the sections below.

5.2 Algorithm for the spectral analysis of the TKE budget

We perform a spectral analysis of the TKE budget following the two-point correlation function approach [121, 84, 88]. Note that there is another equivalent approach that starts from the Navier–Stokes equations in the form of Fourier modes [85, 83]. We adopt the former approach, which is briefly described below. Considering the two-point correlation of $u_i(x, y, z, t)$ and $\tilde{u}_j(x, r_x, y, z, t) = u_j(x + r_x, y, z, t)$, $j = 1, 2, 3$, the spectral

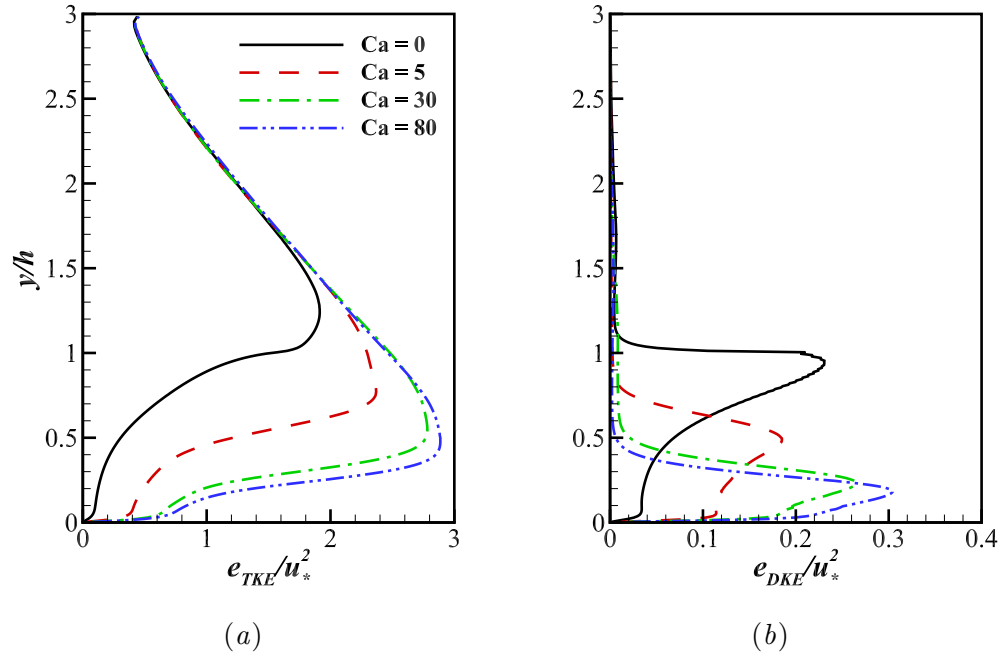


Figure 5.1: Vertical profiles of the (a) TKE and (b) DKE averaged over time and over the x - z planes. Both TKE and DKE are normalized by u_*^2 . The cases are —, $Ca = 0$; - - -, $Ca = 5$; - · - · -, $Ca = 30$; and · · · ·, $Ca = 80$.

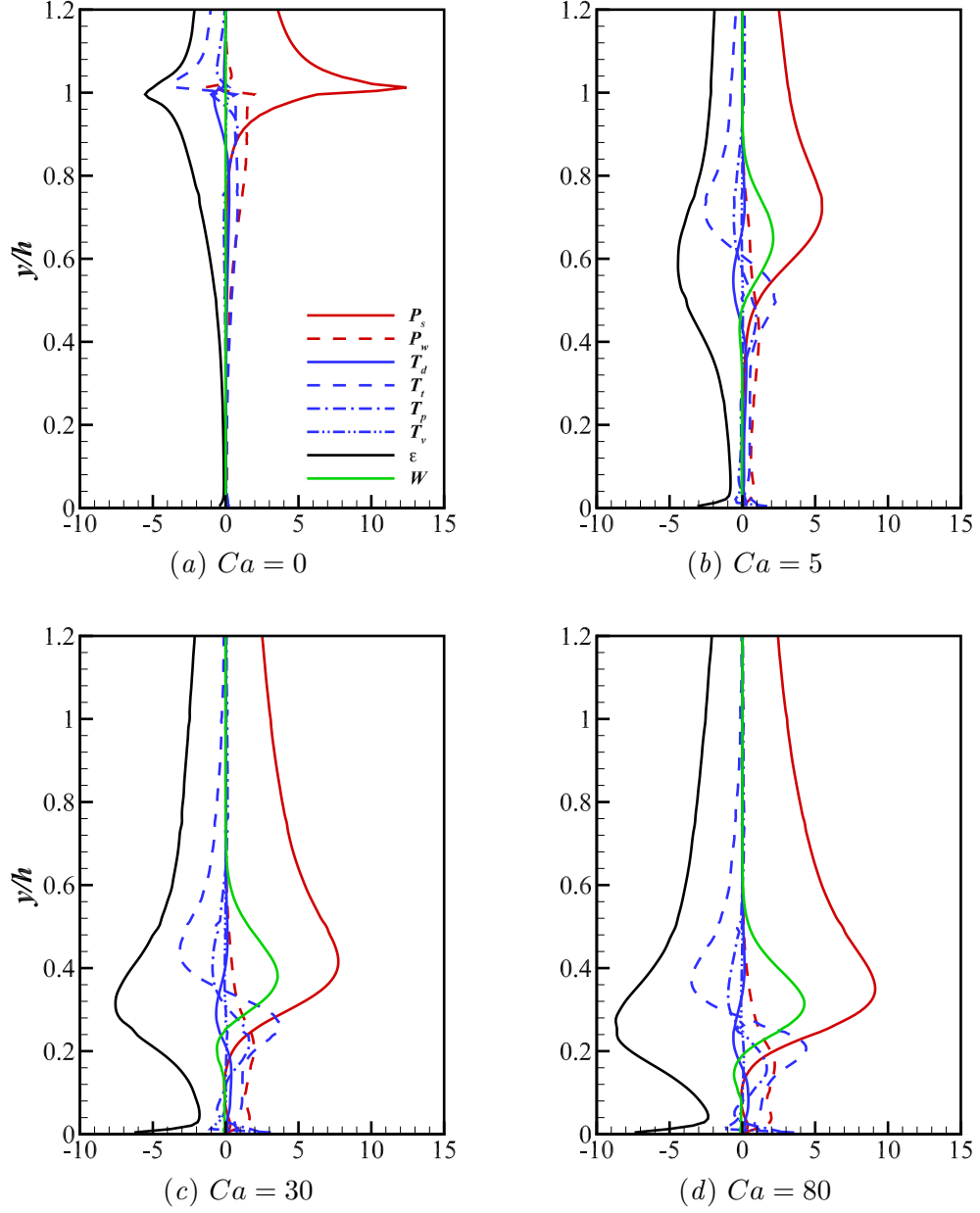


Figure 5.2: Vertical profiles of the TKE budget terms defined in (5.2) inside and near the canopy top for $Ca = 0, 5, 30$, and 80 . The terms are as follows: P_s , the shear production term; P_w , the wake production term; T_d , the dispersive transport term; T_t , the turbulent transport term; T_p , the pressure transport term; T_v , the viscous transport term; ϵ , the dissipation term; and W , the waving term. All the terms are normalized by u^{*3}/h .

budget of the Reynolds stress is

$$\begin{aligned}
\frac{\partial \langle \overline{u'_i \tilde{u}'_j} \rangle_x}{\partial t}(r_x, y, z) = & \underbrace{\left(-\langle \tilde{u}'_j u'_k \rangle_x \frac{\partial \langle \bar{u}_i \rangle_x}{\partial x_k} - \langle u'_i \tilde{u}'_k \rangle_x \frac{\partial \langle \bar{u}_j \rangle_x}{\partial \tilde{x}_k} \right)}_{R_{ij}^P} - \underbrace{\left\langle \tilde{u}'_j \frac{\partial u'_k u'_i}{\partial x_k} + u'_i \frac{\partial \tilde{u}'_k \tilde{u}'_j}{\partial \tilde{x}_k} \right\rangle_x}_{R_{ij}^T} \\
& - \underbrace{\frac{1}{\rho} \left\langle \tilde{u}'_j \frac{\partial p'}{\partial x_i} + u'_i \frac{\partial \tilde{p}}{\partial \tilde{x}_j} \right\rangle_x}_{R_{ij}^\Pi} + \underbrace{\nu \left\langle \tilde{u}'_j \frac{\partial^2 u'_i}{\partial x_k \partial x_k} + u'_i \frac{\partial^2 \tilde{u}'_j}{\partial \tilde{x}_k \partial \tilde{x}_k} \right\rangle_x}_{R_{ij}^\nu} + \underbrace{\langle \tilde{u}'_j f'_i + u'_i \tilde{f}'_j \rangle_x}_{R_{ij}^W},
\end{aligned} \tag{5.3}$$

where $\langle \cdot \rangle_x$ indicates the average in the streamwise direction. In the above equation, R_{ij}^P , R_{ij}^T , R_{ij}^Π , R_{ij}^ν , and R_{ij}^W are the shear production term, turbulent term, pressure term, viscous term, and waving term, respectively. The terms R_{ij}^T , R_{ij}^Π , and R_{ij}^ν can be further decomposed. R_{ij}^T can be decomposed into a wall-normal turbulent transport term R_{ij}^{NTT} and an inter-scale transport term R_{ij}^{IST} :

$$R_{ij}^{NTT} = -\frac{1}{2} \left(\frac{\partial \langle \tilde{u}'_j u'_i u'_2 \rangle_x}{\partial y} + \frac{\partial \langle u'_i \tilde{u}'_j \tilde{u}'_2 \rangle_x}{\partial y} \right), \tag{5.4}$$

$$\begin{aligned}
R_{ij}^{IST} = & \frac{\partial \langle \tilde{u}'_j u'_i u'_1 \rangle_x}{\partial r_x} - \frac{\partial \langle u'_i \tilde{u}'_j \tilde{u}'_1 \rangle_x}{\partial r_x} + \frac{\partial \langle \tilde{u}'_j u'_i u'_3 \rangle_x}{\partial z} - \frac{\partial \langle u'_i \tilde{u}'_j \tilde{u}'_3 \rangle_x}{\partial z} \\
& - \frac{1}{2} \left(\frac{\partial \langle \tilde{u}'_j u'_i u'_2 \rangle_x}{\partial y} + \frac{\partial \langle u'_i \tilde{u}'_j \tilde{u}'_2 \rangle_x}{\partial y} \right) + \left\langle u'_i u'_2 \frac{\partial \tilde{u}'_j}{\partial y} \right\rangle_x + \left\langle \tilde{u}'_j \tilde{u}'_2 \frac{\partial u'_i}{\partial y} \right\rangle_x.
\end{aligned} \tag{5.5}$$

R_{ij}^Π can be decomposed into a wall-normal pressure transport term R_{ij}^{NPT} and an inter-component transport term R_{ij}^{ICT} :

$$R_{ij}^{NPT} = \frac{1}{\rho} \left(\frac{\partial \langle \tilde{u}'_j \tilde{p}' \rangle_x}{\partial y} \delta_{i2} + \frac{\partial \langle u'_i \tilde{p}' \rangle_x}{\partial y} \delta_{j2} \right), \quad (5.6)$$

$$R_{ij}^{ICT} = -\frac{1}{\rho} \left\langle p' \frac{\partial \tilde{u}'_j}{\partial \tilde{x}_i} + \tilde{p}' \frac{\partial u'_i}{\partial x_j} \right\rangle_x. \quad (5.7)$$

R_{ij}^ν can be decomposed into a wall-normal viscous transport term R_{ij}^{NVT} and a dissipation term R_{ij}^{DIS} :

$$R_{ij}^{NVT} = \nu \frac{\partial^2 \langle u'_i \tilde{u}'_j \rangle_x}{\partial y \partial y}, \quad (5.8)$$

$$R_{ij}^{DIS} = -2\nu \left\langle \frac{\partial u'_i}{\partial r_x} \frac{\partial \tilde{u}'_j}{\partial r_x} + \frac{\partial u'_i}{\partial y} \frac{\partial \tilde{u}'_j}{\partial y} + \frac{\partial u'_i}{\partial z} \frac{\partial \tilde{u}'_j}{\partial z} \right\rangle_x. \quad (5.9)$$

Therefore, we have

$$\frac{\partial \langle u'_i \tilde{u}'_j \rangle_x}{\partial t}(r_x, y, z) = R_{ij}^P + R_{ij}^{NTT} + R_{ij}^{IST} + R_{ij}^{NPT} + R_{ij}^{ICT} + R_{ij}^{NVT} + R_{ij}^{DIS} + R_{ij}^W. \quad (5.10)$$

By taking the trace of this equation ($i = j$), performing the Fourier transform about r_x while applying the Wiener-Khinchin theorem and averaging in the spanwise direction, we can obtain

$$\frac{\partial E_e}{\partial t}(k_x, y) = E_P + E_{NTT} + E_{NPT} + E_{NVT} + E_{IST} + E_{DIS} + E_W, \quad (5.11)$$

where k_x is the streamwise wavenumber, E_e is the streamwise spectrum of TKE, and the RHS terms are the shear production E_P , wall-normal turbulent transport E_{NTT} , wall-normal pressure transport E_{NPT} , wall-normal viscous transport E_{NVT} , inter-scale transport E_{ICT} , dissipation E_{DIS} , and waving effect E_W , respectively. The definitions of these terms are as follows:

$$E_e = \Re \left[\left\langle \overline{\mathcal{F}^*(u'_i) \mathcal{F}(u'_i)} \right\rangle_z \right], \quad (5.12a)$$

$$E_P = \Re \left[-2 \left\langle \overline{\mathcal{F}^*(u'_k \frac{\partial u_i}{\partial x_k}) \mathcal{F}(u'_i)} \right\rangle_z \right], \quad (5.12b)$$

$$E_{NTT} = \Re \left[\frac{\partial}{\partial y} \left\langle \overline{\mathcal{F}^*(u'_i u'_2) \mathcal{F}(u'_i)} \right\rangle_z \right], \quad (5.12c)$$

$$E_{NPT} = \Re \left[-\frac{2}{\rho} \left\langle \overline{\mathcal{F}^*(p') \mathcal{F}(u'_2)} \right\rangle_z \right], \quad (5.12d)$$

$$E_{NVT} = \Re \left[\nu \frac{\partial^2}{\partial y^2} \left\langle \overline{\mathcal{F}^*(u'_i) \mathcal{F}(u'_i)} \right\rangle_z \right], \quad (5.12e)$$

$$E_{IST} = \Re \left[-2ik_x \left\langle \overline{\mathcal{F}^*(u'_i u'_1) \mathcal{F}(u'_i)} \right\rangle_z - \frac{\partial}{\partial y} \left\langle \overline{\mathcal{F}^*(u'_i u'_2) \mathcal{F}(u'_i)} \right\rangle_z \right. \\ \left. + 2 \left\langle \overline{\mathcal{F}^*(u'_i u'_2) \mathcal{F}(\frac{\partial u'_i}{\partial y})} \right\rangle_z + 2 \left\langle \overline{\mathcal{F}^*(u'_i u'_3) \mathcal{F}(\frac{\partial u'_i}{\partial z})} \right\rangle_z \right], \quad (5.12f)$$

$$E_{DIS} = \Re \left[-2\nu k_x^2 \left\langle \overline{\mathcal{F}^*(u'_i) \mathcal{F}(u'_i)} \right\rangle_z - 2\nu \left\langle \overline{\mathcal{F}^* \left(\frac{\partial u'_i}{\partial y} \right) \mathcal{F} \left(\frac{\partial u'_i}{\partial y} \right)} \right\rangle_z \right. \\ \left. - 2\nu \left\langle \overline{\mathcal{F}^* \left(\frac{\partial u'_i}{\partial z} \right) \mathcal{F} \left(\frac{\partial u'_i}{\partial z} \right)} \right\rangle_z \right], \quad (5.12g)$$

$$E_W = \Re \left[-2 \left\langle \overline{\mathcal{F}^*(f'_i) \mathcal{F}(u'_i)} \right\rangle_z \right], \quad (5.12h)$$

where $\langle \cdot \rangle_z$ represents the spanwise average, $\mathcal{F}$ is the Fourier transform operator, $(\cdot)^*$ is the complex conjugate operator, and $\Re(\cdot)$ takes the real part of a complex number.

Note that E_{ICT} does not exist in (5.11) because it transfers energy only among the three velocity components. The three components of $\mathbf{E}_{ICT}$ are

$$E_{ICT,i} = \Re \left[\frac{2}{\rho} \left\langle \overline{\mathcal{F}^*(p') \mathcal{F} \left(\frac{\partial u'_i}{\partial x_i} \right)} \right\rangle_z \right], \text{ no summation, } i = 1, 2, \text{ or } 3. \quad (5.13)$$

Here, we consider only the streamwise Fourier modes because we focus on the interaction between the canopy and the streamwise structure of the coherent vortices in the mixing layer. Spanwise Fourier modes need to be considered when studying the spanwise structure, which is important for other types of flows [see, e.g., 83] but is less important than the dominant streamwise structures in canopy flows. In the following analyses of the TKE budget terms, we compare our results mainly with the slow-wave case ($c/u^* = 2$) in [88], who analysed the spectral TKE budget for a three-dimensional turbulent airflow over water waves and has the closest simulation setup in the literature with our canopy problem for spectral TKE budget analysis. Note that because the velocity fluctuations in [88] contain both turbulent and wave-coherent components, which are equivalent to the combination of the turbulent and dispersive components in our case, their TKE budget is equivalent to the sum of the TKE and DKE budgets in our case. Because the DKE budget is nontrivial only at the wake scale, our TKE budget can still be compared with their results, but not at the wake scale. In the following sections § 5.3–5.6, we first discuss the TKE budget in the spectral space at various heights, followed by a discussion of the vertical profiles of the TKE and DKE budgets at the *monami* and wake scales in

§ 5.7.

5.3 Spectra of TKE and the production term

Figure 5.3 shows the streamwise spectra of TKE, E_e . The TKE is higher at low wavenumbers and over the canopy and achieves the largest value near the canopy top. As the canopy flexibility increases, the TKE maximum value increases, and the height of the maximum value decreases as the canopy height decreases. These variations are consistent with figure 5.1(a). From the spectra, we can identify two significant spatial scales that do not exist in the boundary layer spectrum [62]. The first scale is approximately $k_x h \sim 4\pi$, where the TKE inside the canopy reaches a local maximum value, as indicated in figure 5.3. This scale corresponds to the interval between adjacent stems and hence represents the characteristic scale of the wake behind the stems. We call this scale the ‘wake scale’ with a wavenumber k_w . The local maximum TKE is consistent with the spectral shortcut mechanism [79, 61]. It should be noted that such a local maximum TKE also exists at the integer multipliers of the wake scale, also known as the harmonics (not plotted in the figures), and its amplitude decreases as the multiplier increases. These harmonics can also be observed in the results of [88], who showed that the streamwise spectrum of the TKE for air turbulence over a slow water wave exhibits high TKE at the dominant wavelength scale and harmonics at higher wavenumbers. The second significant scale is approximately $k_x h \sim 0.6$. For clarity, figure 5.4 plots locally magnified views of figure 5.3 to highlight the low wavenumber range of $k_x h \leq 5$

in the lower part of the channel ($y/h \leq 1$). For every case, close to the channel bottom wall, there exists a local maximum TKE near $k_x h \sim 0.6$; this is also the global maximum except in the case of $Ca = 80$, where the spatially averaged component at $k_x h \sim 0$ is higher. Note that $k_x h \sim 0.6$ corresponds to the energetic wavenumber component of the *monami*, as shown in figure 4.3, at which the coherent flow structure near the canopy top influences the lower canopy. We call this scale the ‘*monami* scale’ with a wavenumber k_m . The wake and *monami* scales represent the characteristic large flow scale of the coherent structures in the mixing layer and the small flow scale of the wakes behind the stems, respectively. In the following sections, the flow features at these two scales and the interaction between them are analysed and discussed.

To further understand the source of TKE at different streamwise wavenumbers, the streamwise spectra of the shear production term E_P are shown in figure 5.5. Similar to the spectra for the TKE, E_P attains its largest value near the canopy top where the shear is strong and is high at low wavenumbers, especially at the *monami* scale [61]. Local maximum values also exist at the wake scale, but the heights of the regions featuring prominent values differ among the cases: for the rigid canopy, there is only one region near the canopy top ($y/h \sim 1$), whereas for the flexible canopies, there is another region near the channel bottom. We believe that this discrepancy is caused by the different heights of the primary shear. For the rigid canopy ($Ca = 0$), the velocity difference across the canopy top is high and the velocity inside the canopy is low, so the primary shear occurs at the canopy top. As the canopy becomes more flexible, the

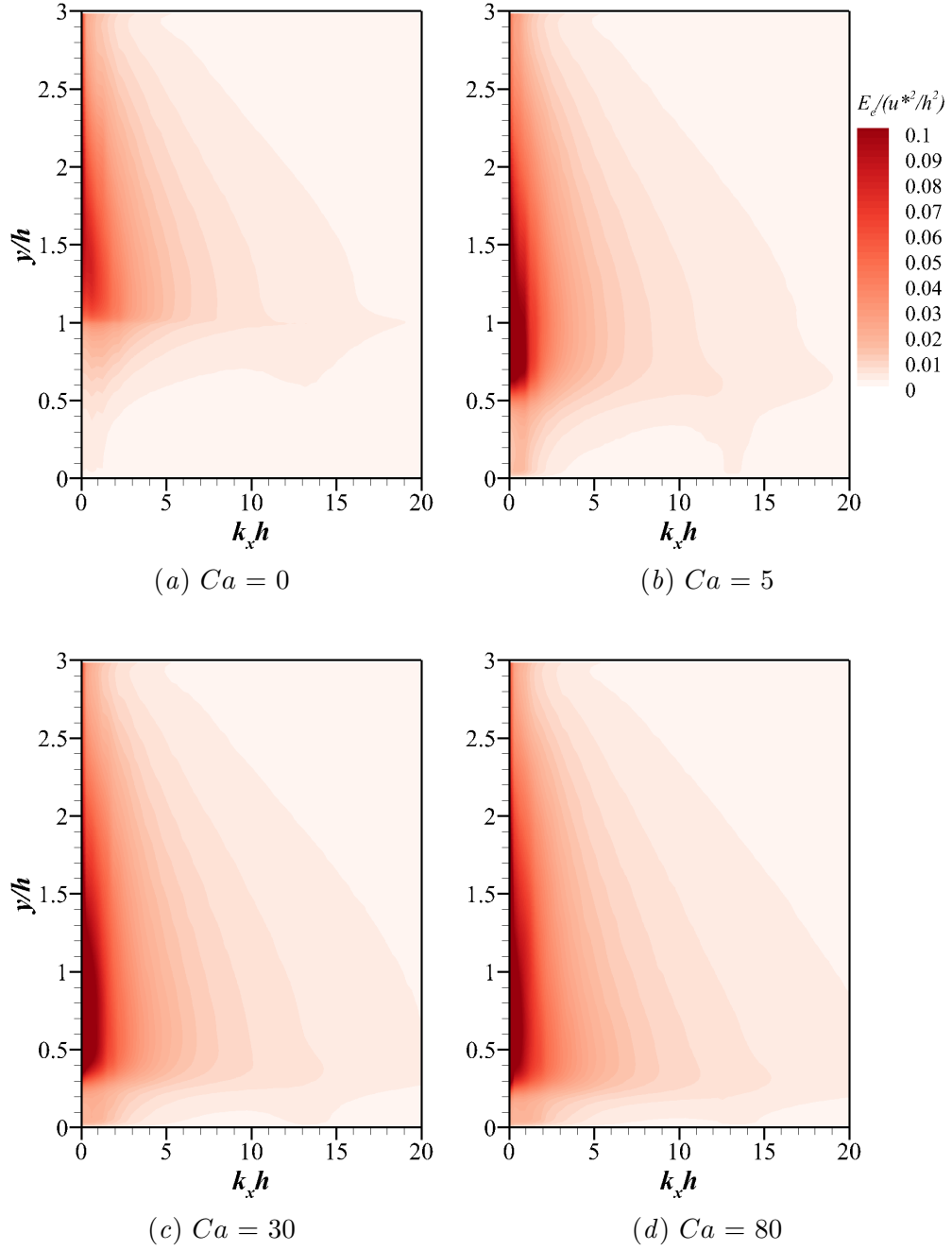


Figure 5.3: One-dimensional streamwise spectra of TKE, E_e , for $Ca = 0, 5, 30$, and 80 . The term is normalized by u^{*2}/h^2 .

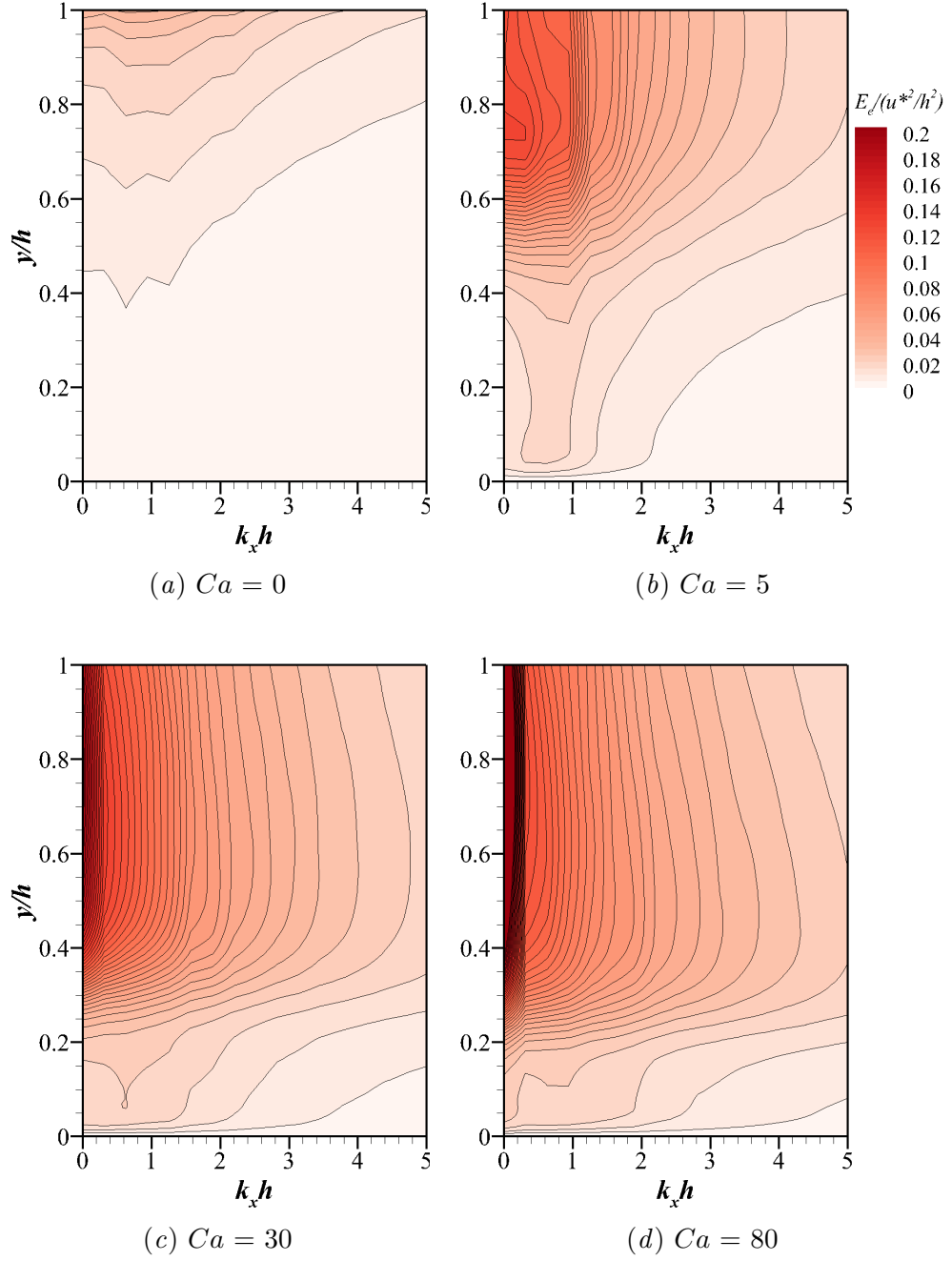


Figure 5.4: Magnified view of figure 5.3 for $k_x h \leq 5$ and $y/h \leq 1$. The increment is 0.005 between contour lines from 0 to 0.2.

boundary between the flows inside and outside the canopy becomes less defined because the oscillation of the stems weakens the shear at the canopy top. At the same time, the flow velocity inside the canopy increases as the drag coefficient decreases owing to the reconfiguration of the stems, and thus, the shear at the channel bottom is strong. The shear strength is indicated by the vertical gradient of the streamwise flow velocity $\partial\langle\bar{u}\rangle/\partial y$ and is observable from figure 4.2(a), where $Ca = 0$ has a high streamwise velocity gradient at the canopy top and $Ca = 30$ and 80 have a high streamwise velocity gradient close to the channel bottom wall.

Note that some previous studies found regions with negative shear production in spectral analyses of the TKE budget; these regions are thought to be related to a negative Reynolds shear stress $-\langle u'v' \rangle$ at certain heights and wavenumbers [84, 88]. Specifically, [88] found a region with negative shear production at the dominant wavelength scale in the case of wind over a slow wave and attributed it to the strong negative Reynolds shear stress inside the critical layer. In our cases, we did not observe any regions with prominent negative shear production, and figure 4.10 excludes a prominent positive correlation between u' and v' at the *monami* scale. The reason is that, different from the case of wind over a slow wave studied by [88], the canopy flow considered herein is predominantly influenced by the coherent structures in the mixing layer, which interact with the flow inside the canopy through sweep and ejection events [116, 72, 117].

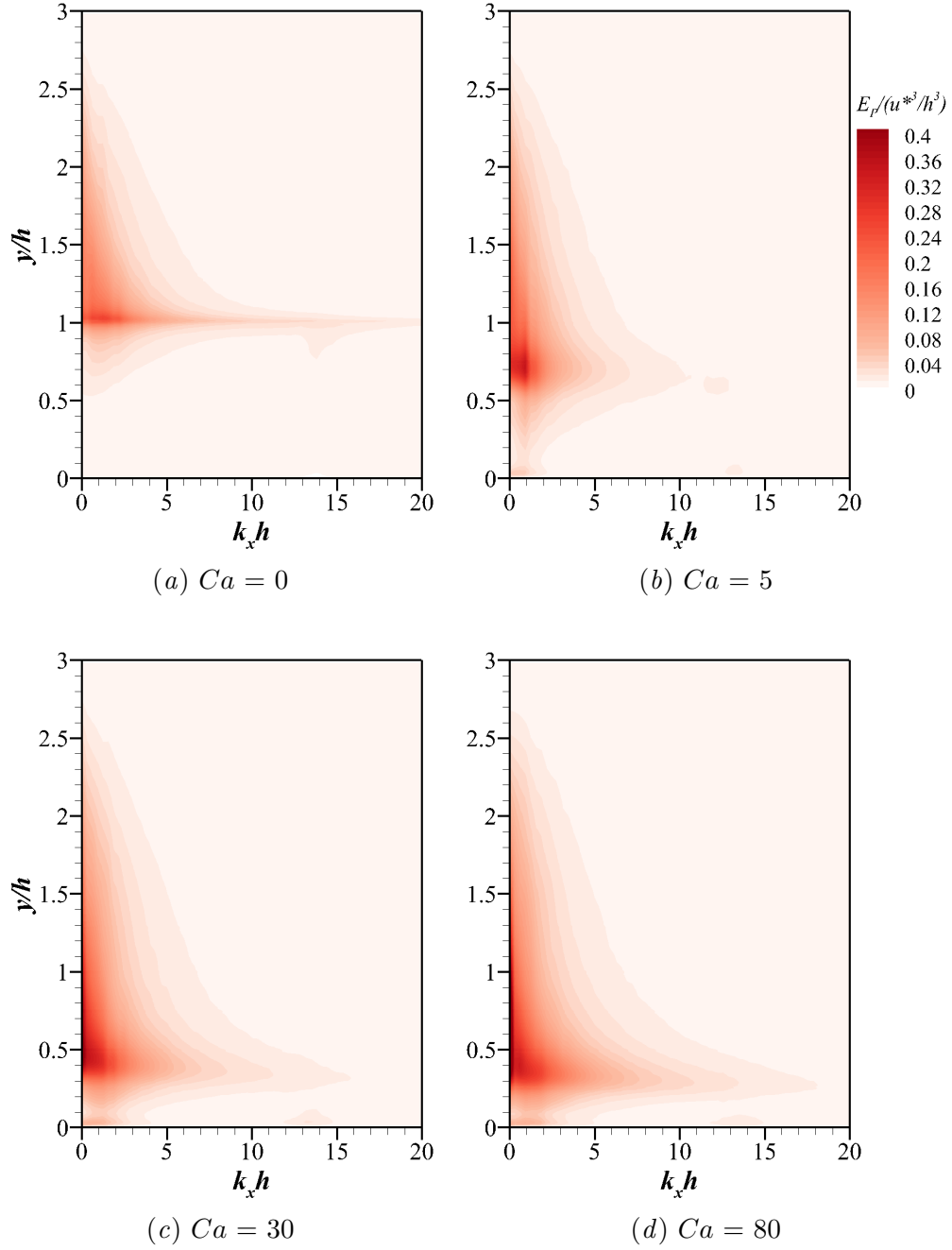


Figure 5.5: One-dimensional streamwise spectra of the production term E_P for $Ca = 0, 5, 30$, and 80 . The term is normalized by u^{*3}/h^3 .

5.4 Wall-normal transport terms

The wall-normal transport terms, including the wall-normal pressure transport term E_{NPT} , the wall-normal turbulent transport term E_{NTT} , and the wall-normal viscous transport term E_{NVT} , transport TKE along the vertical direction. Note that

$$\int_0^H E_{NTT}(k_x, y) dy = \int_0^H E_{NPT}(k_x, y) dy = \int_0^H E_{NVT}(k_x, y) dy = 0. \quad (5.14)$$

Figure 5.6 shows the three wall-normal transport terms. All the cases show similar features near the canopy top, where all three terms mainly transport the TKE downward across the canopy top to balance the high energy production outside the canopy with the high dissipation inside the canopy, similar to the observation from figure 5.2. The amplitudes of the pressure transport term E_{NPT} and the turbulent transport term E_{NTT} increase as the canopy flexibility increases. For the viscous transport term E_{NVT} , its amplitude and variation along the vertical direction are larger at the top of the rigid canopy than at the tops of the flexible canopies because the shear is sharper at the former but is smeared by the oscillation of stems in the flexible canopies. The amplitudes of the three terms are different. Comparing the ranges of the amplitudes for these terms in figure 5.6 indicates that E_{NTT} has the largest amplitude, followed by E_{NPT} , and E_{NVT} has the smallest amplitude. In [88], although the turbulent transport term also achieves the largest amplitude, the amplitude of the pressure transport term is smaller than that of the viscous transport term. We believe that this discrepancy is related to the relative

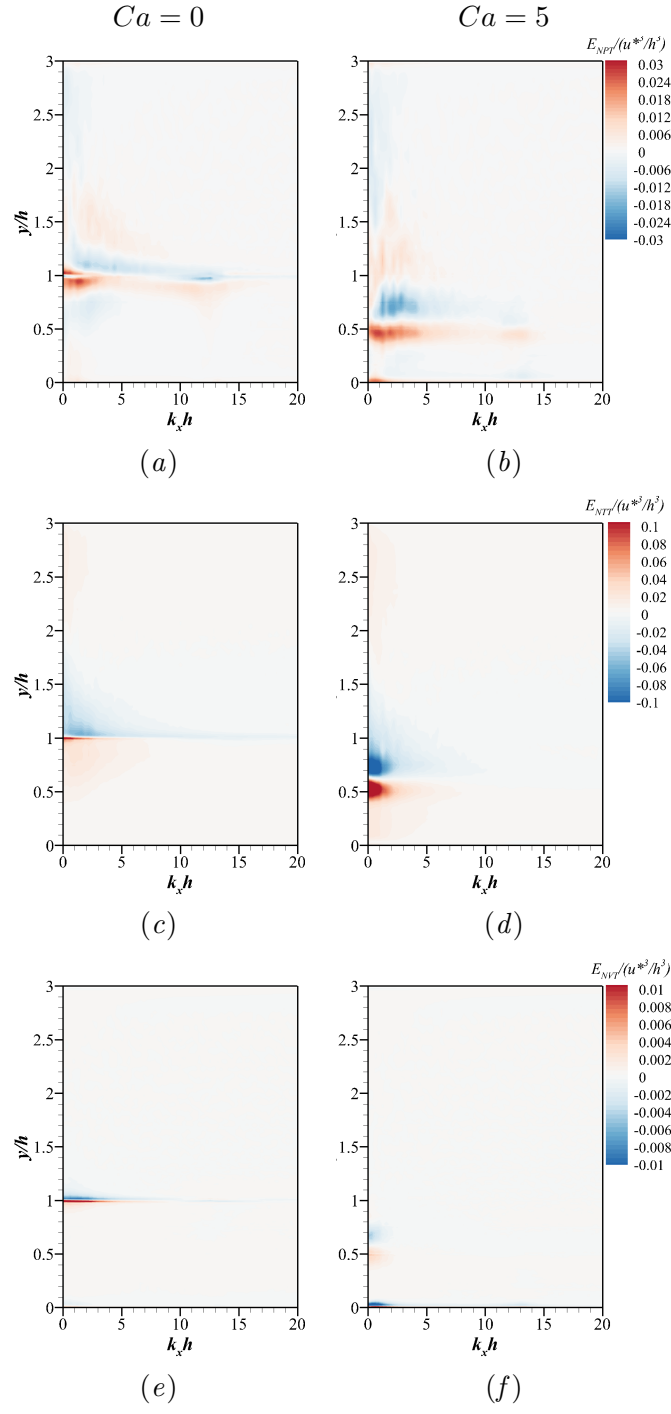
amplitude of the pressure drag and viscous drag on the wave surface. In [88], because the wave amplitude is small, the pressure drag is small, and thus, the viscous drag is the dominant drag on the wave surface, so the wall-normal transport associated with the viscous effect is high. In our cases, however, the canopy drag is the predominant form of drag, which results in high wall-normal transport associated with the pressure variation. In addition to the different absolute amplitudes, the amplitude attenuation rate with increasing wavenumber also varies among E_{NTT} , E_{NPT} and E_{NVT} , as shown in figure 5.6. In particular, the pressure transport term has a slower rate than the other two terms, so its effects on large and small scales only slightly differs.

5.5 Inter-component and inter-scale terms

The inter-component terms $E_{ICT,i}(i = 1, 2, 3)$ transport TKE among three velocity components such that

$$E_{ICT,1}(k_x, y) + E_{ICT,2}(k_x, y) + E_{ICT,3}(k_x, y) = 0. \quad (5.15)$$

These terms are plotted in figure 5.7. All the cases show similar results except for the vertical locations of their variations owing to the different canopy heights. Outside the canopy, the TKE generally flows from the streamwise component to the vertical and spanwise components, consistent with the conclusions of previous spectral TKE analyses of channel flows [121, 84, 85, 83]. However, at scales larger than the wake scale



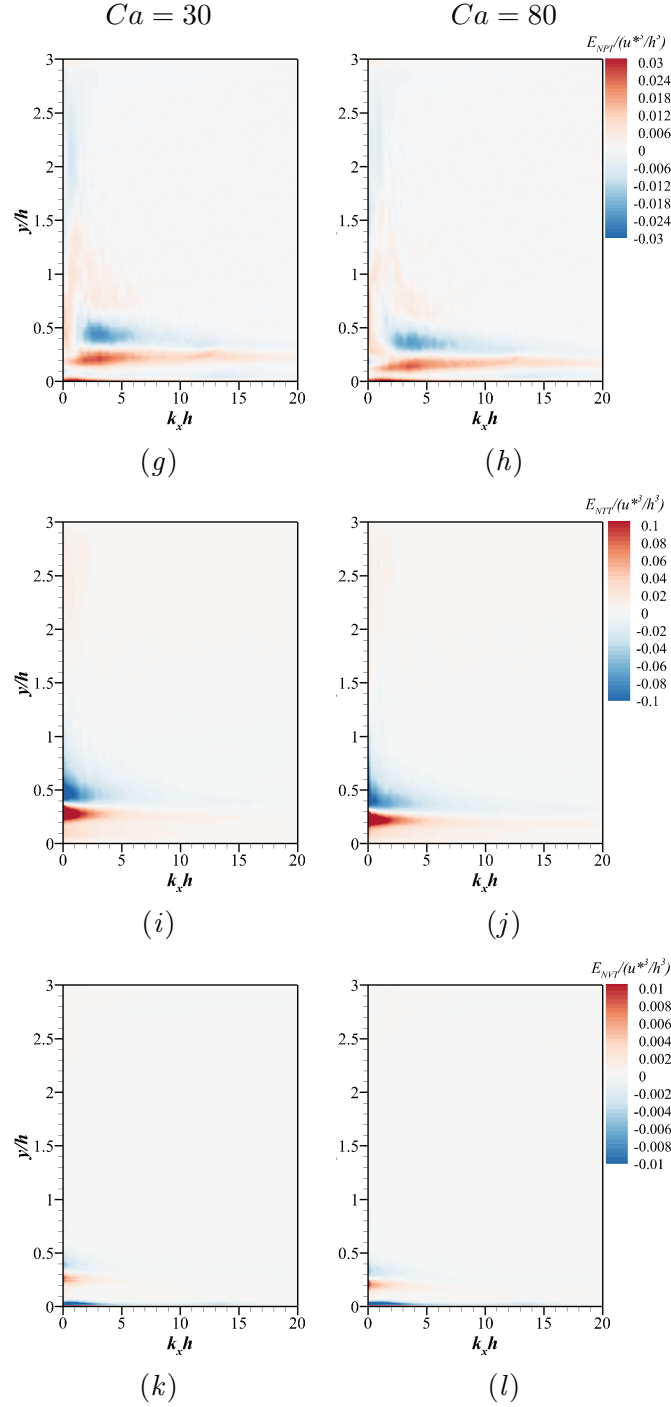


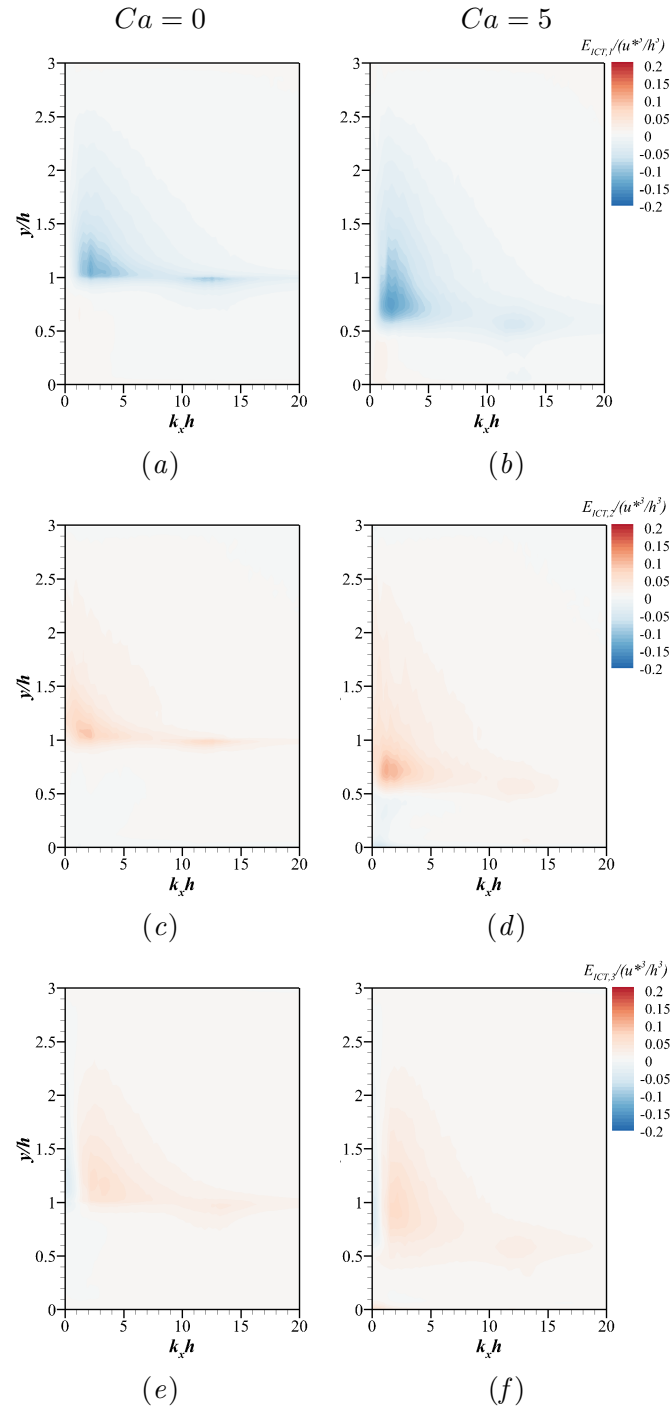
Figure 5.6: One-dimensional streamwise spectra of the (a,b,g,h) wall-normal pressure transport term E_{NPT} , (c,d,i,j) wall-normal turbulent transport term E_{NTT} , and (e,f,k,l) wall-normal viscous transport term E_{NVT} . The terms are normalized by u^*3/h^3 . The cases are (a,c,e) $Ca = 0$, (b,d,f) $Ca = 5$, (g,i,k) $Ca = 30$, and (h,j,l) $Ca = 80$.

($k_x < k_w$), the TKE flows from the spanwise component to the vertical component. A similar phenomenon was also observed in previous studies [121, 84, 88]. [84] pointed out that the source of the TKE in the spanwise component is the inter-scale transport, and we observe a similar phenomenon. The inter-component transport inside the canopy is generally small, and only for small wavenumbers ($k_x h < 5$) is there a weak TKE flow from the vertical component to the streamwise and spanwise components. The reason for the above phenomenon is that inside the canopy, high hydrodynamic pressure ($p' > 0$) usually exists on the windward side of the stems with large deformation, thereby driving or sustaining their deformation, and these stems correspond to large deformation events and large forward speed events (§ 4.4), both of which correspond to sweep events in which the negative vertical velocity increases with height ($\partial v' / \partial y < 0$), as shown in figure 4.10(*b, d*).

The inter-scale transport term E_{IST} transports TKE across different scales such that

$$\int_0^\infty E_{IST}(k_x, y) dk_x = 0. \quad (5.16)$$

The inter-scale transport term is shown in figure 5.8, where the isopleth of $E_{IST} = 0$ is plotted to show the boundary between the TKE recipient and donor regions. For all the cases, this term extracts TKE from the low-wavenumber range with $k_x h \lesssim 5$ as well as the scales near the wake scale k_w inside the canopy. Note that the maximum resolvable wavenumber in our study is $k_{\max} h = 40\pi$, which is far larger than the wavenumbers at



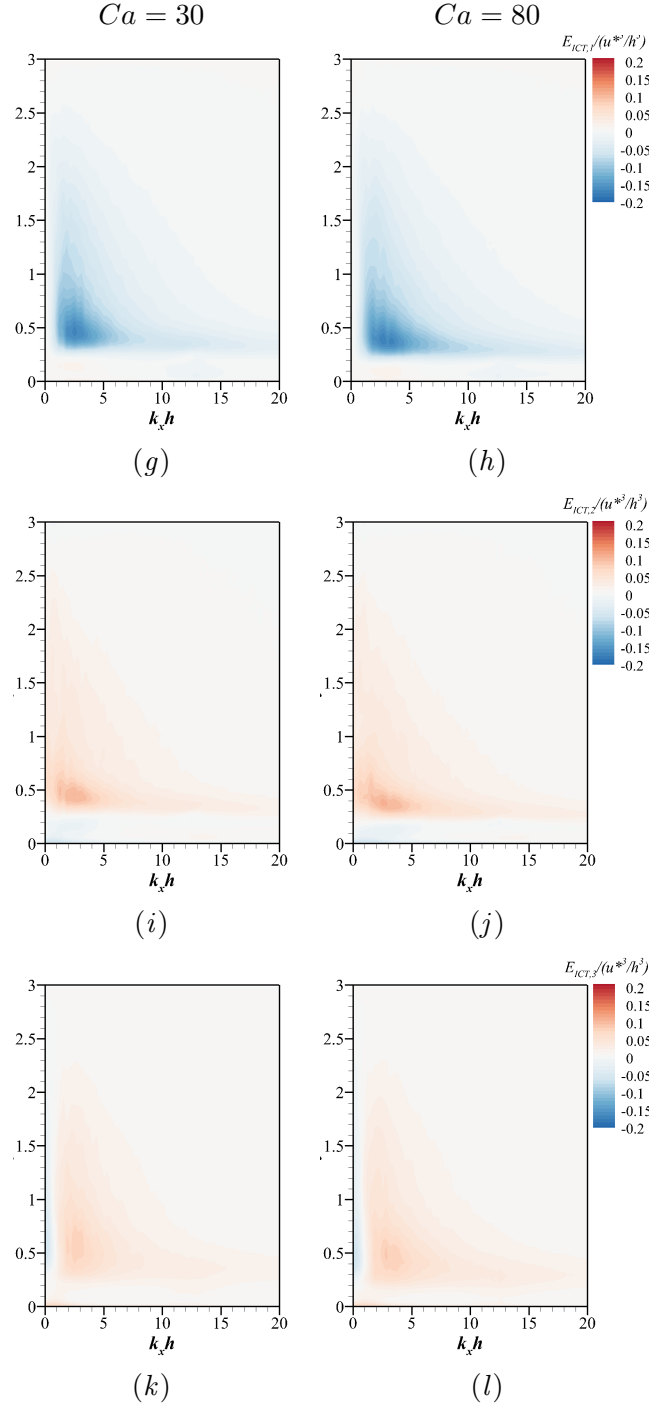


Figure 5.7: One-dimensional streamwise spectra of the inter-component transport terms of different velocity components: (a,b,g,h) the streamwise component $E_{ICT,1}$; (c,d,i,j) the vertical component $E_{ICT,2}$; and (e,f,k,l) the spanwise component $E_{ICT,3}$. The terms are normalized by u^*/h^3 . The cases are (a,c,e) $Ca = 0$, (b,d,f) $Ca = 5$, (g,i,k) $Ca = 30$, and (h,j,l) $Ca = 80$.

which TKE is extracted by this term. Therefore, the TKE extracted at small wavenumbers is distributed across a broad wavenumber range such that the positive values of E_{IST} do not appear significant in the figure compared with the negative values. We checked that (5.16) holds for our data. Note that while the spectral shortcut mechanism may suggest an inter-scale energy flow route from the shear production scale to the wake scale, according to figure 5.8, TKE is extracted by the inter-scale term from the wake scale inside the canopy. Therefore, the spectral shortcut mechanism is unlikely to be exerted by the inter-scale transport term. We find that it is instead associated with the waving term, which is discussed in § 5.6.

The energy cascade from large to small scales is ubiquitous in turbulent flows. Previous studies on channel flow also demonstrated the existence of an inverse energy cascade through which the TKE is transferred from small to large scales [121, 84, 83, 122, 88]. In our cases, as shown in figure 5.8, there is a TKE recipient region near the bottom wall with a wavenumber range of $5 \lesssim k_x h \lesssim 10$ between two donor regions. Therefore, we examine which donor region is the TKE source of the recipient region to understand whether the inverse energy cascade occurs in our case. For this purpose, we perform a triadic analysis of the inter-scale transport term E_{IST} . Note that previous studies [e.g., 83, 88] performed triadic analyses on the turbulent transport term E_T , which is the sum of the wall-normal turbulent transport term E_{NTT} and the inter-scale transport term E_{IST} :

$$E_T = E_{NTT} + E_{IST} = -\Re \left[\left\langle \widehat{u_i'}^* \frac{\partial}{\partial x_j} \widehat{u_i' u_j'} \right\rangle_z \right], \quad (5.17)$$

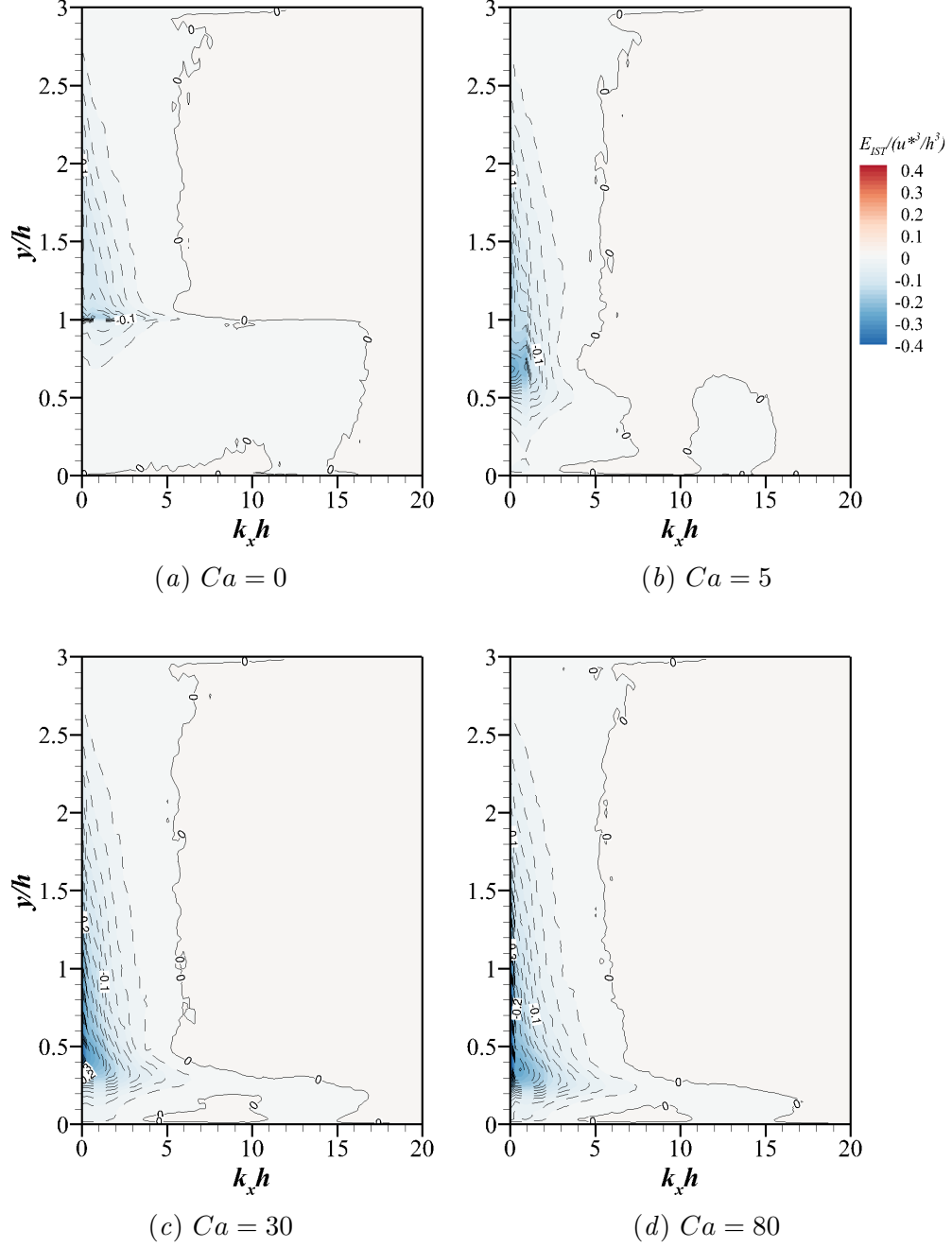


Figure 5.8: One-dimensional streamwise spectra of the inter-scale transport term E_{IST} normalized by u^{*3}/h^3 for $Ca = 0, 5, 30$, and 80 . The isopleths of the normalized E_{IST} are also plotted, where the solid lines represent nonnegative values of $E_{IST}/(u^{*3}/h^3) = 0$ and the dashed lines represent negative values with an interval of 0.02 .

where $\widehat{(\cdot)}$ is the Fourier transform operator. According to (5.12) and (5.17), the formula for E_{IST} is complex, while the formulas for E_T and E_{NTT} are simple. Therefore, to perform the triadic analysis of E_{IST} , we first perform triadic analyses of E_T and E_{NTT} and then subtract the result of E_{NTT} from that of E_T . In what follows, we take E_T as an example to show the triadic analysis process. According to the convolution theorem, we have

$$\begin{aligned}
\widehat{u'_i u'_j}(k_0) &= \int_{k_l+k_m=k_0} \widehat{u'_i}(k_l) \widehat{u'_i}(k_m) dk_l = \int_{\substack{k_l+k_m=k_0 \\ k_l, k_m \geq 0}} \widehat{u'_i}(k_l) \widehat{u'_i}(k_m) dk_l \\
&+ \int_{\substack{k_m-k_l=k_0 \\ k_l, k_m \geq 0}} \widehat{u'_i}(k_l) \widehat{u'_i}(k_m) dk_l + \int_{\substack{k_l-k_m=k_0 \\ k_l, k_m \geq 0}} \widehat{u'_i}(k_l) \widehat{u'_i}(k_m) dk_l.
\end{aligned} \tag{5.18}$$

In this equation, the last three terms indicate the excitation at wavenumber k_0 by the interaction of wavenumbers k_l and k_m . The term with $k_l + k_m = k_0$ indicates the interaction of two large scales for a small scale and thus represents a normal energy cascade. The terms with $k_m - k_l = k_0$ and $k_l - k_m = k_0$ indicate the interaction of a large scale and a small scale for an intermediate scale and thus represent an inverse

energy cascade. Therefore, for E_T , we have

$$\begin{aligned}
 E_T(k_0, y) &= \int_{k_0 - k_{\max}}^{k_{\max}} E_{T,c}(k_0, y, k_l) dk_l \\
 &= \int_{k_0 - k_{\max}}^{k_{\max}} -\Re \left\langle \widehat{u}_i^*(k_0, y, z, t) \frac{\partial}{\partial x_j} \left[\widehat{u}_i'(k_l, y, z, t) \widehat{u}_j'(k_0 - k_l, y, z, t) \right] \right\rangle_z dk_l.
 \end{aligned} \tag{5.19}$$

Then, by performing the same triadic analysis on E_{NTT} for $E_{NTT,c}$ and subtracting the result for $E_{NTT,c}$ from that for $E_{T,c}$, we obtain

$$\begin{aligned}
 E_{IST}(k_0, y) &= \int_{k_0 - k_{\max}}^{k_{\max}} E_{IST,c}(k_0, y, k_l) dk_l \\
 &= \int_{k_0 - k_{\max}}^{k_{\max}} E_{T,c}(k_0, y, k_l) dk_l - \int_{k_0 - k_{\max}}^{k_{\max}} E_{NTT,c}(k_0, y, k_l) dk_l,
 \end{aligned} \tag{5.20}$$

where $E_{NTT,c}$ is obtained in a way similar to (5.19):

$$E_{NTT,c}(k_0, y, k_l) = -\Re \left[\frac{\partial}{\partial y} \left\langle \widehat{u}_i'(k_l, y, z, t) \widehat{u}_2'(k_0 - k_l, y, z, t) \widehat{u}_i^*(k_0, y, z, t) \right\rangle_z \right]. \tag{5.21}$$

Considering the symmetry of wavenumbers k_l and $k_0 - k_l$ in the triadic interaction with wavenumber k_0 , we further average $E_{IST,c}(k_0, y, k_l)$ and $E_{IST,c}(k_0, y, k_0 - k_l)$ for the final result. Figure 5.9 shows the inter-scale turbulent energy transport to the wake scale k_w via triadic interaction, $[E_{IST,c}(k_w, y, k_l) + E_{IST,c}(k_w, y, k_w - k_l)]/2$. Over the canopy, its value is positive for $0 < k_l < k_w$, which reflects the energy cascade as the energy flows

to the wake scale from the large scale. Inside the canopy, the energy influx to the wake scale k_w originates from the interaction between the mean flow $k_l = 0$ and the wake scale itself, not from the interactions among different turbulence scales. Additionally, inside the canopy, for $0 < k_l h < k_w h$, the value of the triadic interaction term is negative but not significant, which suggests that the outflux from the wake scale to the large scale is weak if it does exist. Therefore, we conclude that the inverse cascade is insignificant in our cases.

5.6 Waving term

Following (3.20) and (5.1), the waving term for the total energy $W_{total} = \langle \overline{f_i u_i} \rangle$ can be decomposed into

$$W_{total} = \underbrace{\langle \overline{f_i} \rangle \langle \overline{u_i} \rangle}_{W_{MKE}} + \underbrace{\langle \overline{f_i''} \overline{u_i''} \rangle}_{W_{DKE}} + \underbrace{\langle \overline{f_i'} \overline{u_i'} \rangle}_{W_{TKE}}, \quad (5.22)$$

where the RHS terms are the waving terms for MKE, DKE, and TKE.

Figure 5.10 shows the vertical profiles of W_{MKE} , W_{DKE} , W_{TKE} , and W_{total} . The waving term for MKE is a sink because $\langle \overline{f_i} \rangle < 0$ and $\langle \overline{u_i} \rangle > 0$. That for DKE is a source because in the dominant streamwise component (see figure 5.11), the negative dispersive velocity $\overline{u_1''} < 0$ and negative dispersive stem drag $\overline{f_1''} < 0$ are correlated on the stems and $\overline{u_1''} > 0$ and $\overline{f_1''} > 0$ are correlated at the interval between stems. For the rigid canopy case, as the waving terms for TKE and the total energy are zero (§ 5.1), the waving effect purely transfers energy from MKE to DKE. In contrast, for the flexible

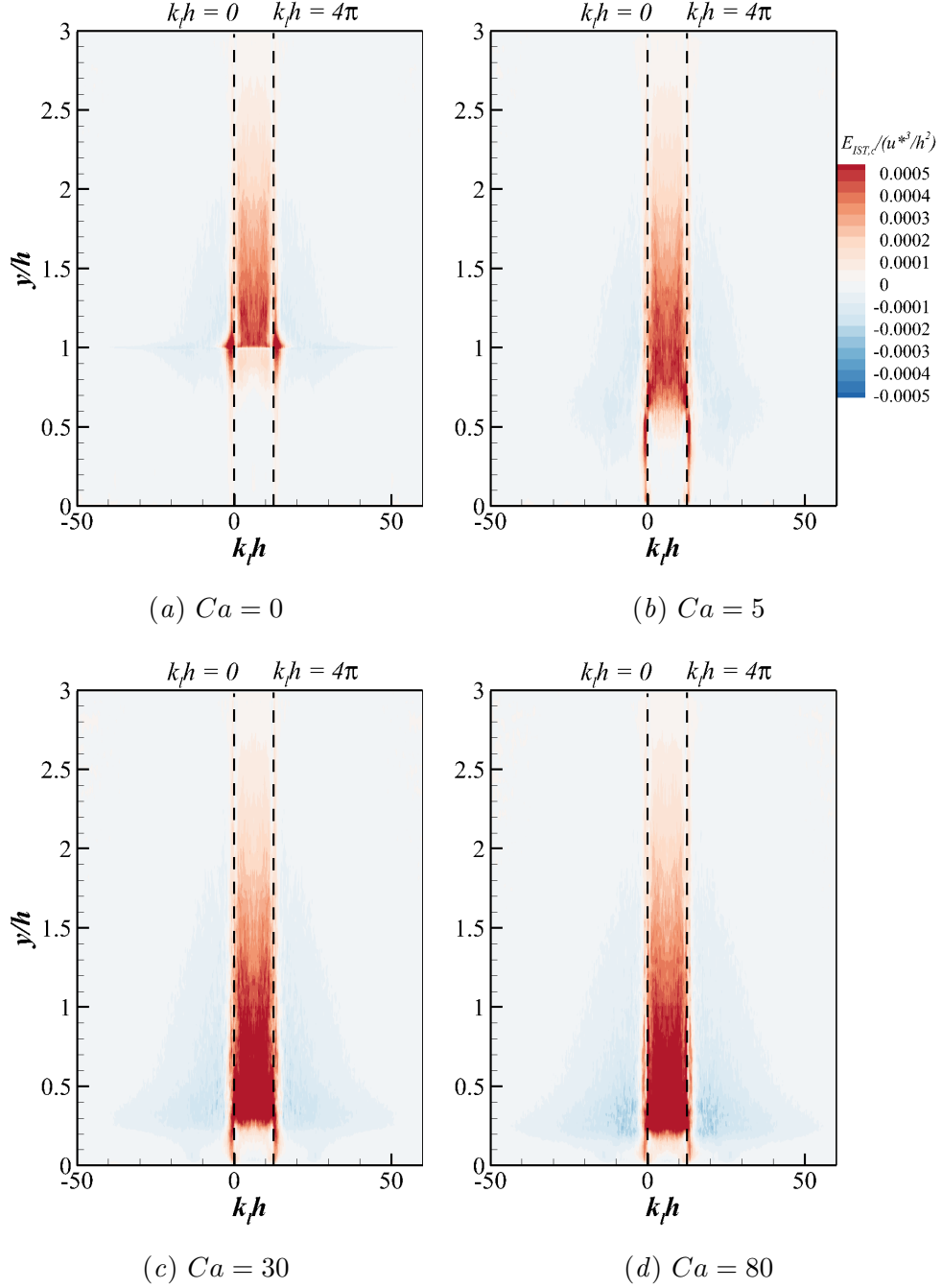


Figure 5.9: Normalized inter-scale turbulent energy transport to the wake scale k_w via triadic interaction, $[E_{IST,c}(k_w, y, k_l) + E_{IST,c}(k_w, y, k_w - k_l)]/2/(u^{*3}/h^2)$, defined by (5.20) for cases $Ca = 0, 5, 30$, and 80 . The two vertical dashed lines represent the average flow scale $k_l h = 0$ and the wake scale $k_l h = k_w h = 4\pi$.

canopy cases, the waving term for TKE is an energy source near the canopy top and a weak energy sink inside the canopy, as discussed in § 5.1. Additionally, the effect of the waving term on the total energy is an energy sink, which means that it dissipates energy. As shown in figure 5.10, both W_{MKE} and W_{DKE} have two local extrema: one near the canopy top and one near the bottom wall. The upper extremum is due to the strong flow and large drag near the canopy top, while the lower extremum is correlated with the mean velocity variation shown in § 4.2.

To better understand the effects of these waving terms, we analyse the three components W_{MKE} , W_{DKE} , and W_{TKE} , as shown in figure 5.11. Only the streamwise component contributes to W_{MKE} because $\langle \bar{u}_2 \rangle = \langle \bar{u}_3 \rangle = 0$. For W_{DKE} , the contribution predominantly comes from the streamwise component, except in the flexible canopy case with $Ca = 5$, where the vertical component $W_{DKE,2}$ provides a comparable contribution in the upper canopy. This is because for both the rigid canopy ($Ca = 0$) and the highly flexible canopies ($Ca = 30$ and 80), the stems exhibit no or very small oscillations (see the discussion on ‘high flexibility-induced rigidity’ in § 4.3) such that $\bar{u}_2 \approx 0$ and $\bar{u}_2'' \approx 0$ on the stems, while in the case with $Ca = 5$, the large stem oscillations in the upper canopy stimulate the inhomogeneity of the vertical velocity and drag along the trajectory of the stem oscillation. A similar explanation also applies to the decreasing contribution of $W_{DKE,2}$ as the canopy flexibility increases (figure 5.11*f–h*) because highly flexible stems are more likely to yield to the flow and display smaller oscillation amplitudes compared with the $Ca = 5$ case, as discussed in § 4.3. For W_{TKE} , the

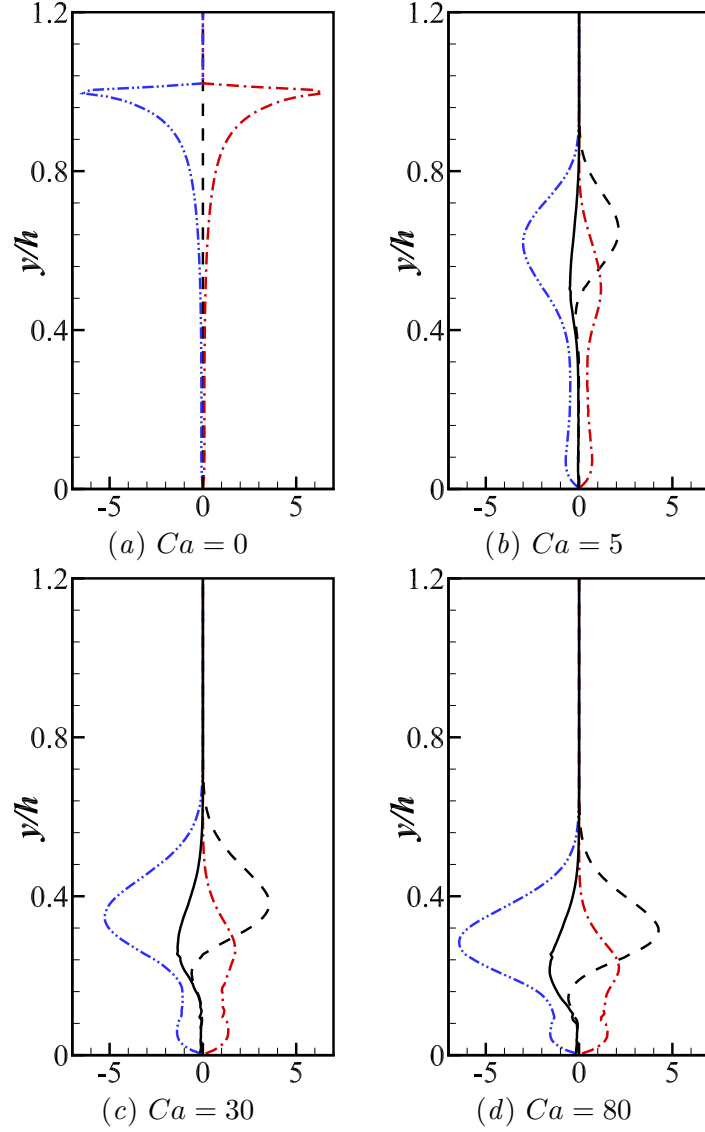
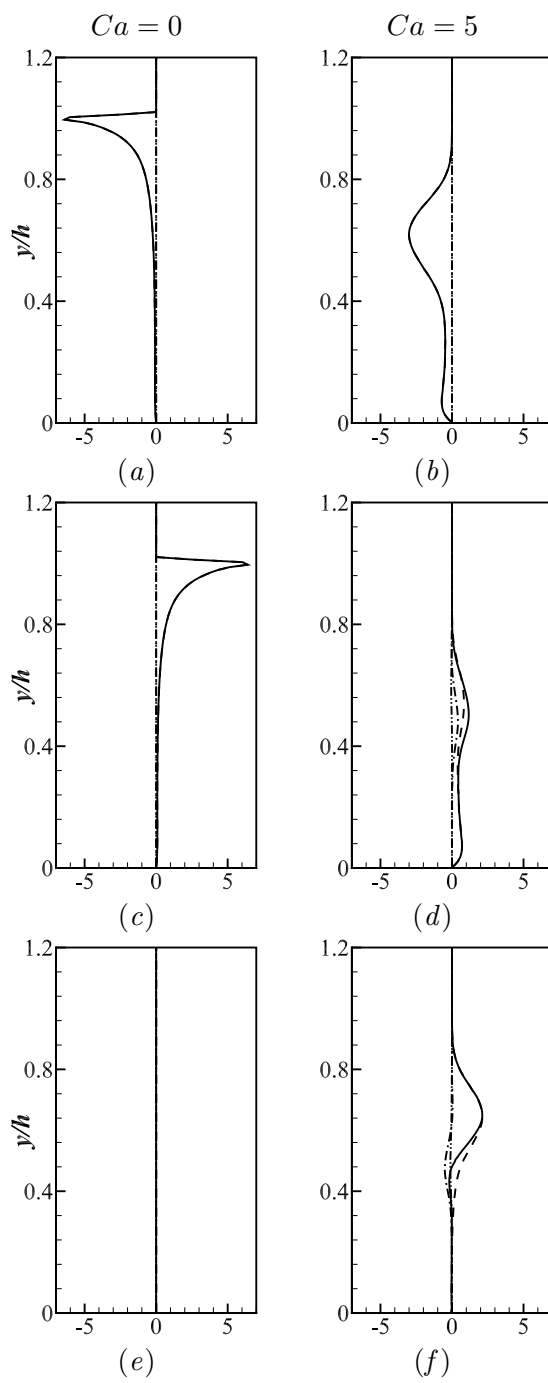


Figure 5.10: Vertical profiles of the waving terms for MKE (W_{MKE}), DKE (W_{DKE}), TKE (W_{TKE}) and the total energy (W_{total}) normalized by u^{*2} for cases $Ca = 0, 5, 30$, and 80 . The lines are as follows: $\text{---}\cdot\text{---}$, W_{MKE} ; $\text{---}\cdot\text{---}$, W_{DKE} ; $\text{---}\text{---}\text{---}$, W_{TKE} ; and --- , W_{total} . Note that $W_{total} = 0$ for $Ca = 0$ and is not shown on (a) to avoid obscuring W_{TKE} .

streamwise and vertical components are a strong source and a weak sink, respectively, and the spanwise component changes from a weak sink to a weak source as the height increases. Therefore, the positive sign of W_{TKE} results from the streamwise component.

To understand the contributions of the different $f'_i-u'_i$ events, we perform a quadrant analysis on each component of the TKE waving term, $\langle \overline{f'_i u'_i} \rangle$ ($i = 1, 2$, or 3 ; no summation notation). The quadrant events for the $f'_i-u'_i$ correlation are defined as Q1 events if $f'_i > 0$ and $u'_i > 0$; Q2 events if $f'_i < 0$ and $u'_i > 0$; Q3 events if $f'_i < 0$ and $u'_i < 0$; and Q4 events if $f'_i > 0$ and $u'_i < 0$. Their contributions to the TKE waving term in the flexible canopy cases are plotted in figure 5.12. All the cases show similar results except for the vertical locations of their variations owing to the different canopy heights. For the streamwise and vertical components, the contributions from Q3 and Q4 events are larger than those from Q1 and Q2 events, respectively. In other words, the events with $u'_i < 0$ contribute more to W_{TKE} than those with $u'_i > 0$. For the streamwise component, this is because $u'_1 < 0$ means that the stem has a larger velocity relative to the streamwise flow such that the drag force is higher. For the vertical component, this is because sweep events (where $u'_2 < 0$) are stronger than ejection events (where $u'_2 > 0$) in the mixing layer [116]. For the spanwise component, the contributions from Q1 and Q2 events are the same as those from Q4 and Q3 events, respectively, which is due to the symmetry in the spanwise direction. The spanwise component is negative because the stems do not deform in the spanwise direction (§ 3.3); hence, the stem drag and spanwise flow velocity component act in opposite directions. As the spanwise



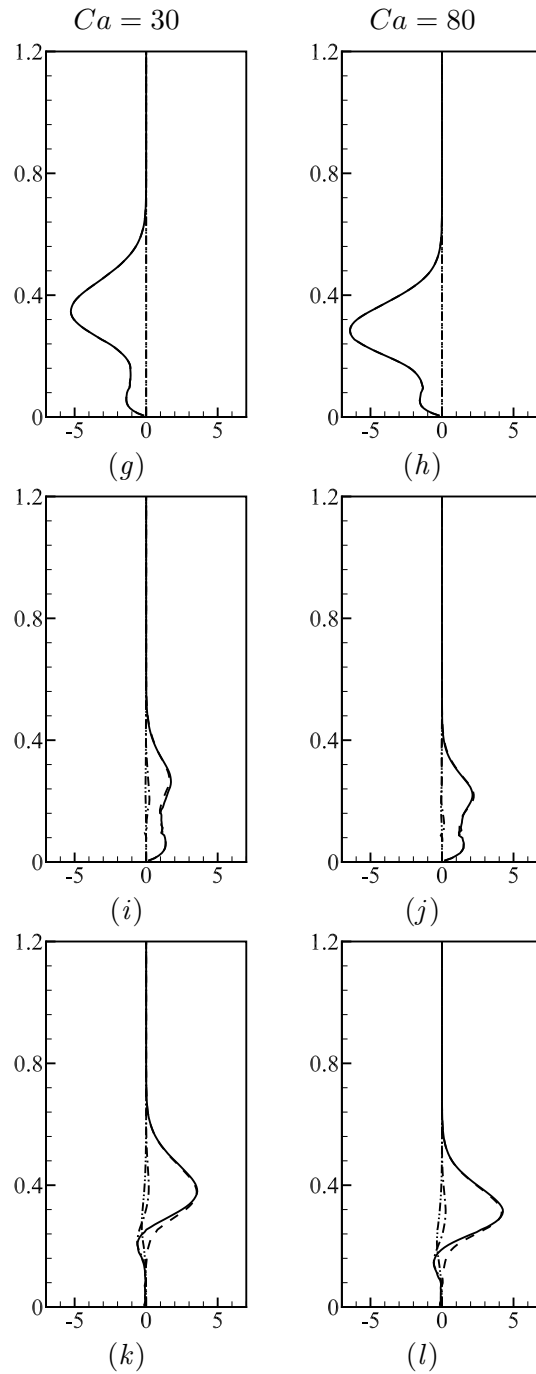


Figure 5.11: Vertical profiles of the waving term for the different velocity components of (a,b,g,h) MKE, (c,d,i,j) DKE, and (e,f,k,l) TKE. The lines are as follows: - - -, the streamwise component; — · —, the vertical component; · · ·, the spanwise component; and —, the sum of all the components. The cases are (a,c,e) $Ca = 0$, (b,d,f) $Ca = 5$, (g,i,k) $Ca = 30$, and (h,j,l) $Ca = 80$.

component has smaller values than the other two components, it is not analysed in the remainder of this section.

We further examine the contributions of different conditionally averaged events to the streamwise and vertical components of the TKE waving term W_{TKE} in the flexible canopy cases. Figure 5.13 shows the contours of the conditionally averaged streamwise component of the waving term, $\langle W_1 \rangle^c$, in the x - y plane cutting through the centre of the stem row. We discuss the different Ca cases together because they have similar results, although the vertical locations of the variations are different owing to differences in the deformation of the canopy. For large deformation and large forward speed events, $\langle W_1 \rangle^c$ over the canopy is positive because $f'_1 > 0$ and $u'_1 > 0$, as explained in § 4.4. Therefore, in these situations, the positive values of $\langle W_1 \rangle^c$ are Q1 events of the f'_1 - u'_1 correlation. For small deformation events, at the centre of the domain, the values are positive around the stem tips and negative in the upper canopy. As shown in figures 4.10 and 4.11, this is because $u'_1 < 0$ at both locations as a result of the ejection events of the turbulence, and f'_1 is negative and positive at the stem tips and upper stems, respectively, as explained in § 4.4. Therefore, the positive and negative values are Q3 and Q4 events, respectively, of the f'_1 - u'_1 correlation. This result indicates a wall-normal TKE transport mechanism by stem waving. For large backward speed events, the positive $\langle W_1 \rangle^c$ values near the stem tips are due to the Q3 events of the f'_1 - u'_1 correlation, i.e., $f'_1 < 0$ and $u'_1 < 0$.

Figure 5.14 is similar to figure 5.13 except that it shows the vertical component of the waving term, $\langle W_2 \rangle^c$. Similar to the discussion in figure 5.13, we discuss all the cases

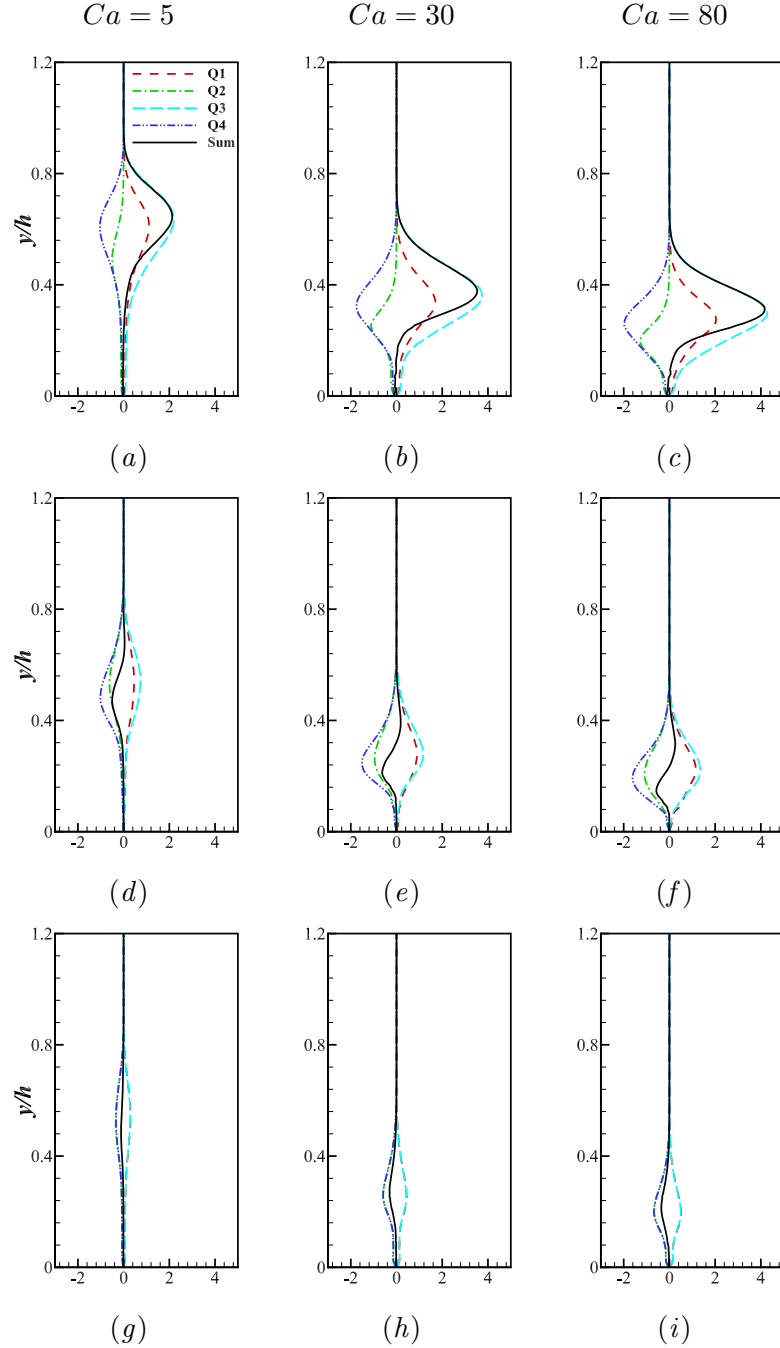
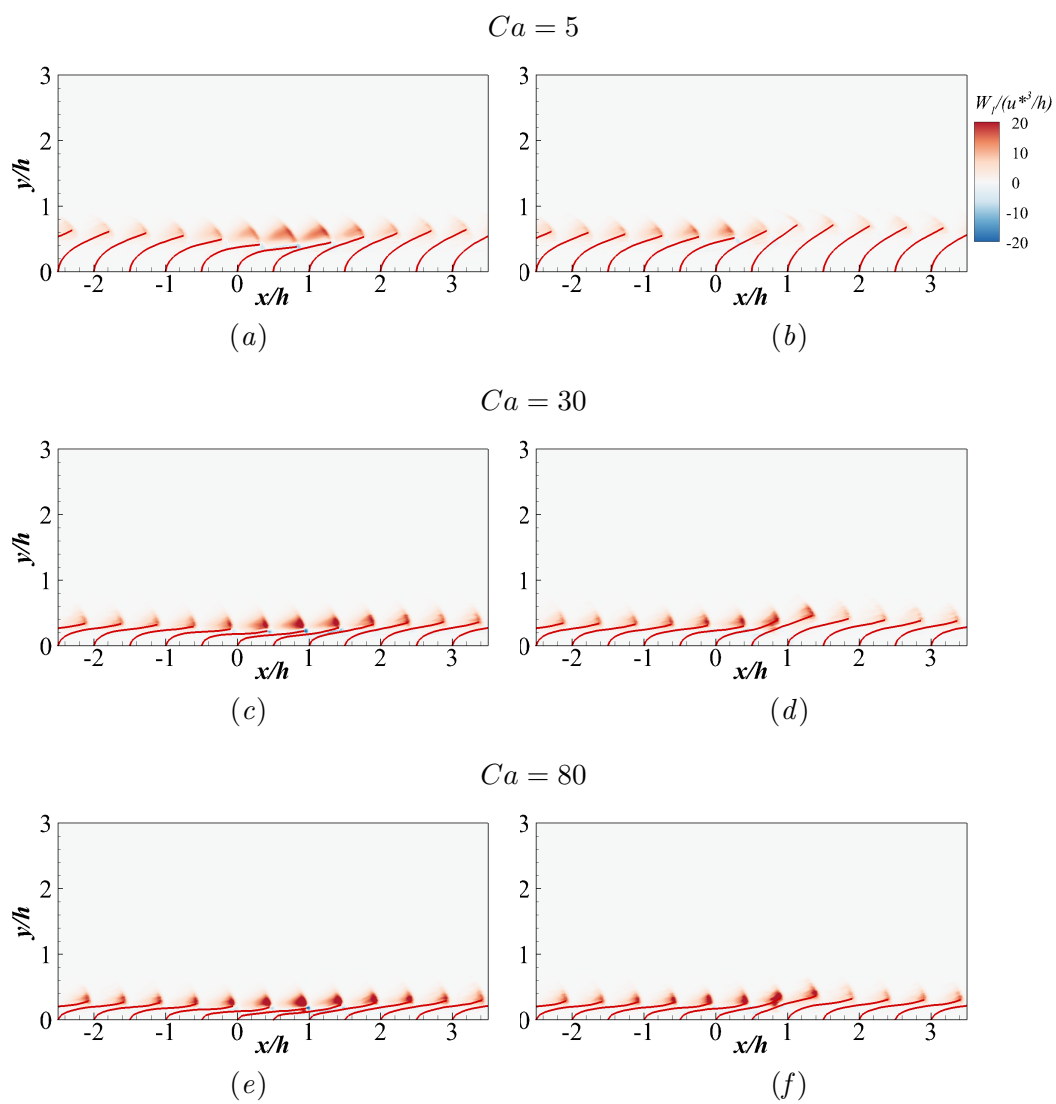


Figure 5.12: Contributions to $f'_i u'_i$ ($i = 1, 2$, or 3 ; no summation notation) from its quadrant events: $---$, Q1 events ($f'_i > 0, u'_i > 0$); $- \cdot -$, Q2 events ($f'_i < 0, u'_i > 0$); $---$, Q3 events ($f'_i < 0, u'_i < 0$); $\cdots$, Q4 events ($f'_i < 0, u'_i > 0$); and $---$, the sum of all the events. The components are (a-c) the streamwise component ($i = 1$), (d-f) the vertical component ($i = 2$), and (g-i) the spanwise component ($i = 3$). The cases are (a,d,g) $Ca = 5$, (b,e,h) $Ca = 30$, and (c,f,i) $Ca = 80$.



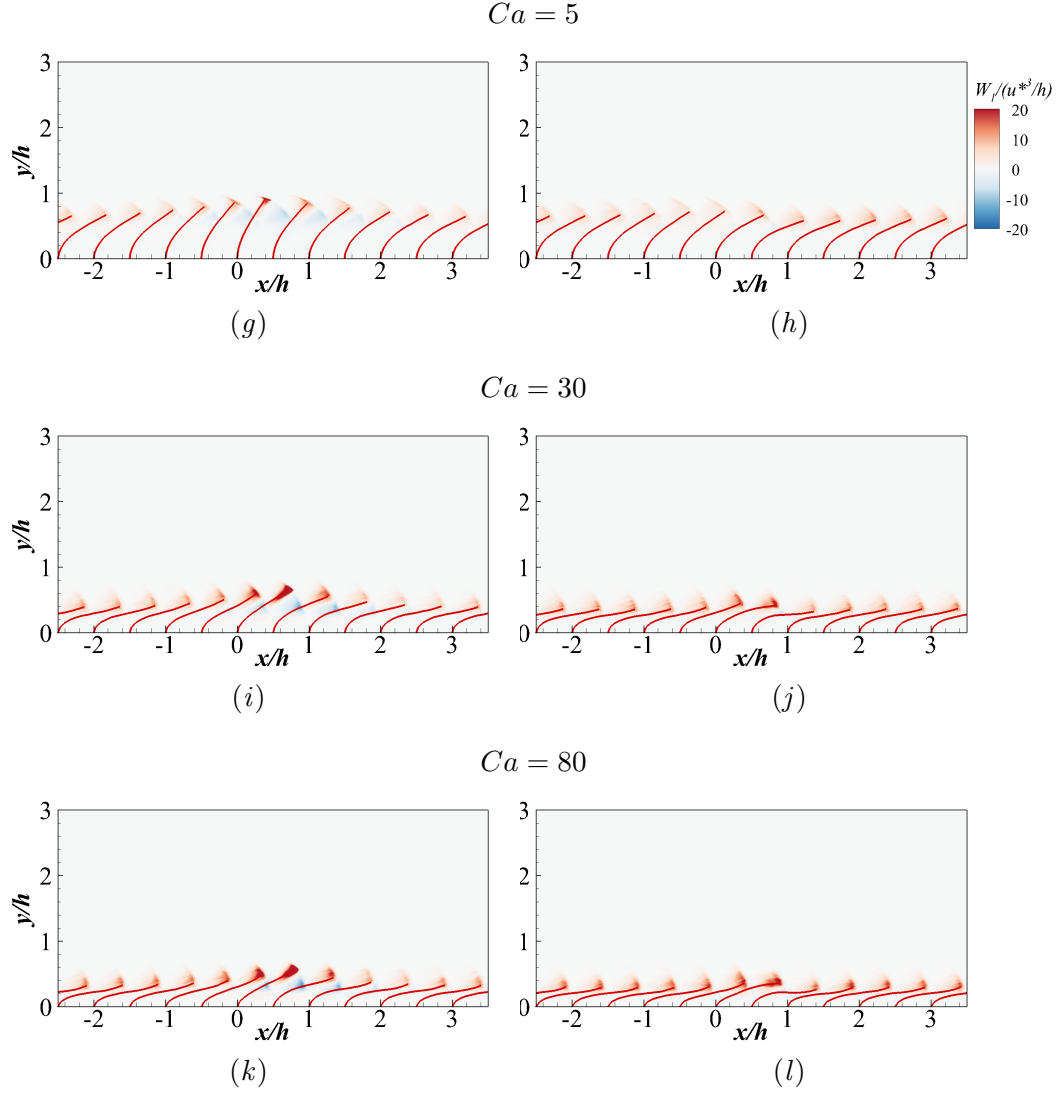
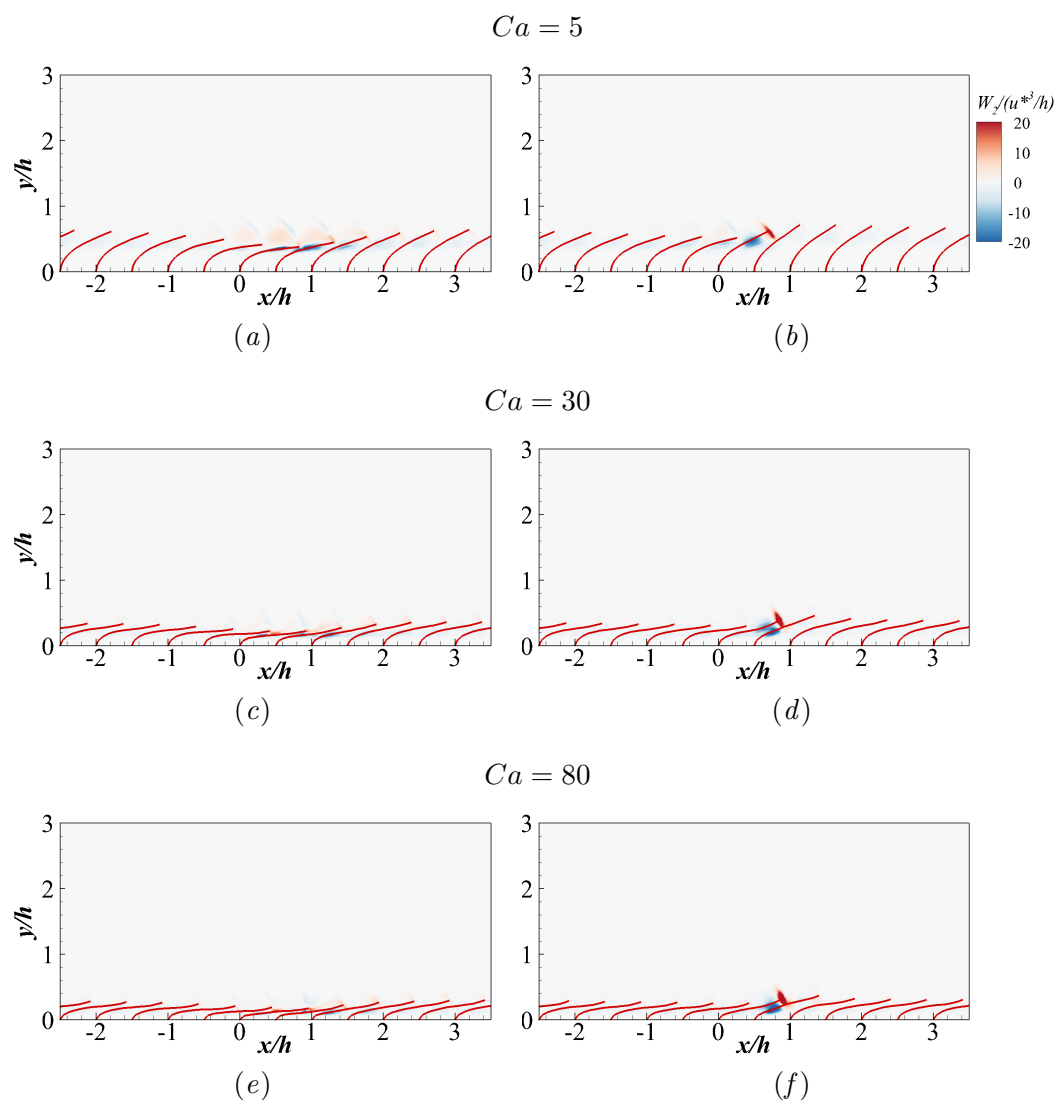


Figure 5.13: Contours of the conditionally averaged normalized streamwise component of the waving term, $\langle W_1 \rangle^c/(u^{*3}/h)$, in the x - y plane cutting through the centre of the stem row. The conditionally averaged events are (a,c,e) large deformation events, (b,d,f) large forward speed events, (g,i,k) small deformation events, and (h,j,l) large backward speed events. The cases are (a,b,g,h) $Ca = 5$, (c,d,i,j) $Ca = 30$, and (e,f,k,l) $Ca = 80$. The red lines represent the conditionally averaged shapes of the stems.

together owing to the similarity of the results. For large deformation and large forward speed events, the positive and negative values near the canopy top are Q3 and Q4 events, respectively, of the $f'_2-u'_2$ correlation because they correspond to sweep events of the turbulent flow where $u'_2 < 0$ (figure 4.10*b,d*). For small deformation and large backward speed events, the positive and negative values in the upper canopy are Q1 and Q2 events, respectively, of the $f'_2-u'_2$ correlation, as they correspond to ejection events of the turbulent flow where $u'_2 > 0$ (figure 4.10*f,h*). Explanations for the sign of f'_2 are given in § 4.4.

Similar to other terms, spectral analysis on the waving term can also be performed to understand the effects of the waving term on energy transport among different vertical locations and scales. The streamwise spectra of the TKE waving term, E_W , are plotted in figure 5.15. For the rigid canopy case, E_W is a sink at low wavenumbers ($k_x h < 10$) and a source at high wavenumbers ($k_x h > 10$). The sink and source effects are strong in the upper canopy and achieve maximal values near the *monami* scale k_m and the wake scale k_w , respectively. These phenomena indicate that E_W plays a considerable role in the inter-scale transport of energy in the rigid canopy case. The flexible canopy cases show similar phenomena to the rigid canopy case at relatively high wavenumbers ($k_x h > 10$). At relatively low wavenumbers ($k_x h < 10$), E_W transports energy upward inside the canopy. Therefore, in flexible canopies, E_W has the combined effects of inter-scale transport and wall-normal transport. For the inter-scale transport effect of E_W , in all cases, the energy flux is from the *monami* scale to the wake scale, which is consistent



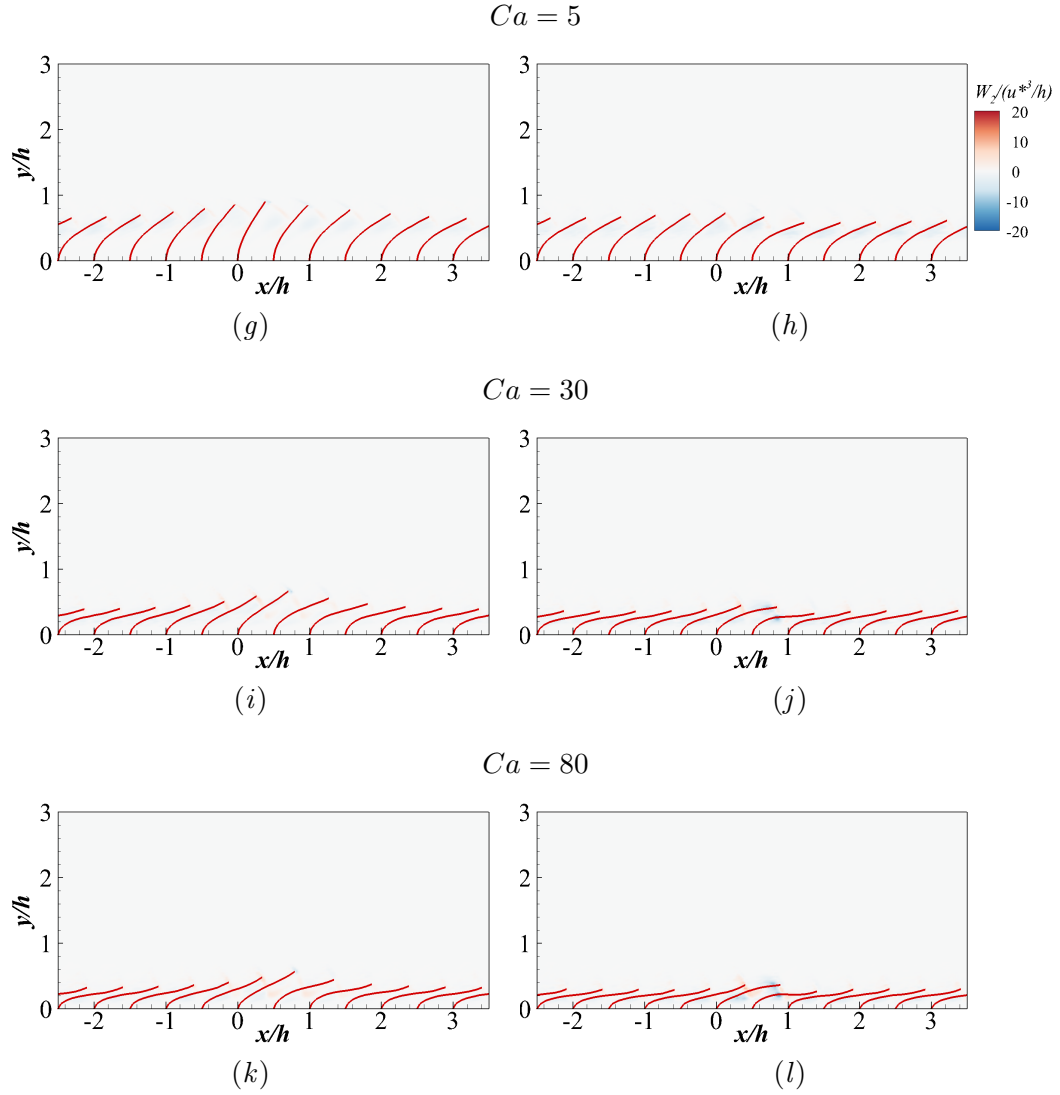


Figure 5.14: Contours of the conditionally averaged normalized vertical component of the waving term, $\langle W_2 \rangle^c/(u^{*3}/h)$, in the x - y plane cutting through the centre of the stem row. The conditionally averaged events are (a,c,e) large deformation events, (b,d,f) large forward speed events, (g,i,k) small deformation events, and (h,j,l) large backward speed events. The cases are (a,b,g,h) $Ca = 5$, (c,d,i,j) $Ca = 30$, and (e,f,k,l) $Ca = 80$. The red lines represent the conditionally averaged shapes of the stems.

with the spectral shortcut mechanism. Therefore, we conclude that the spectral shortcut mechanism is exerted by the waving term associated with the correlation between the drag and velocity fluctuations.

5.7 Vertical profiles of TKE and DKE at the *monami* and wake scales

The *monami* and wake scales are the two characteristic scales in our canopy flow cases. We perform a comparison of the TKE and DKE budgets at these two scales, and the results are plotted in figures 5.16 and 5.17.

Figure 5.16 shows the vertical profiles of the TKE budget terms at the *monami* and wake scales. At the *monami* scale (figure 5.16*a-d*), the globally dominant source and sink are the shear production and inter-scale transport terms, respectively, and the wall-normal turbulent transport term plays a major role in the transport of TKE in the vertical direction. In contrast, the waving term acts as a sink inside the canopy for the rigid canopy and as a transport term for the flexible canopies. In the flexible canopy cases, the signs of the waving and wall-normal turbulent transport terms are opposite to each other and swap at the canopy top. Therefore, the canopy top is the characteristic height that divides the two different patterns of energy transport. Inside the canopy, the shear production and wall-normal turbulent transport terms act as TKE sources at the *monami* scale, while the inter-scale transport and waving terms are sinks. Above

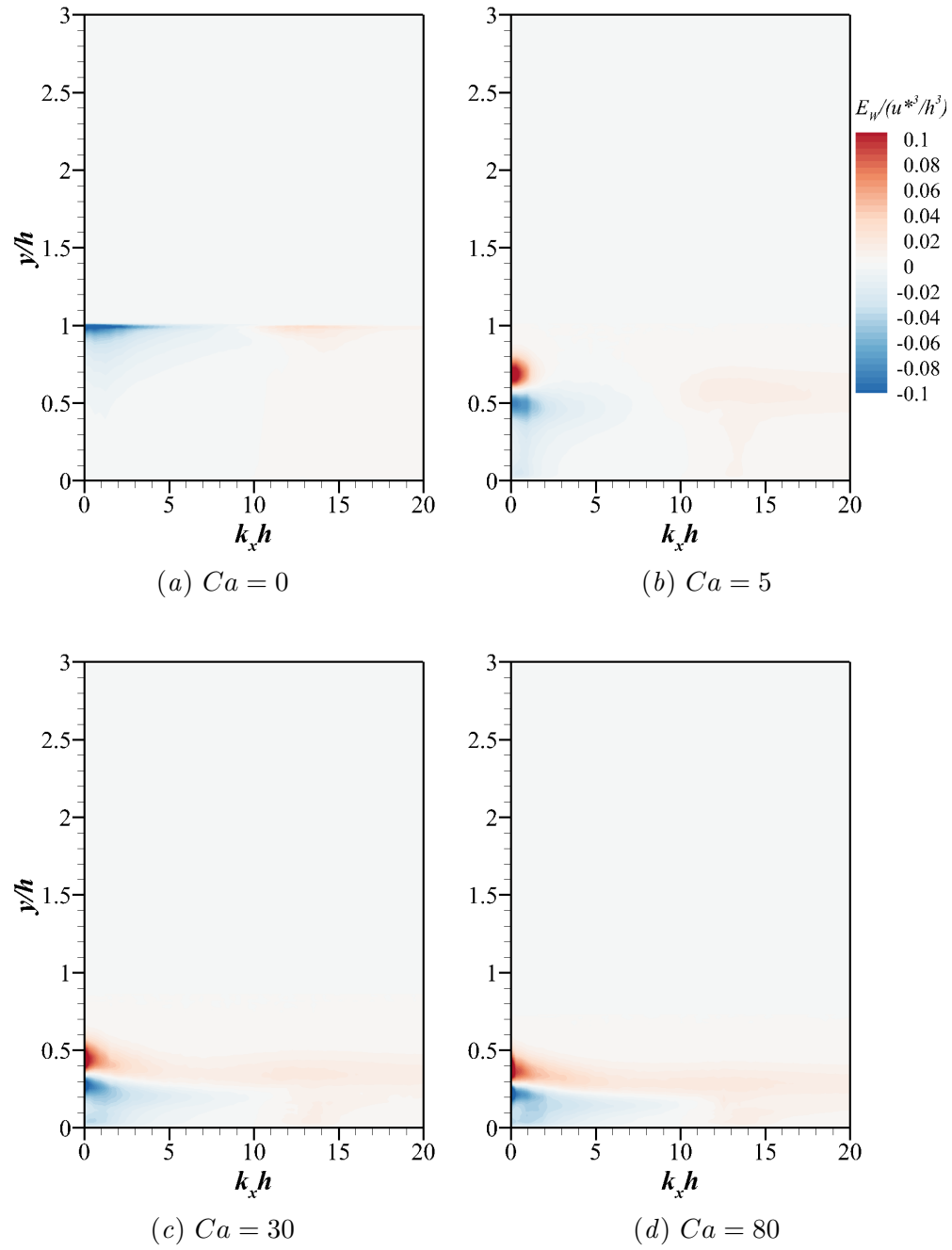


Figure 5.15: One-dimensional streamwise spectra of the waving term E_W for $Ca = 0, 5, 30$, and 80 . The term is normalized by u^{*3}/h^3 .

the canopy, the sources are the shear production and waving terms near the canopy top, and the sinks are the inter-scale transport and wall-normal turbulent transport terms. A summary of the TKE flow at the *monami* scale for the rigid and flexible canopies is given in figure 5.18(*a*) and (*c*), respectively.

Likewise, the TKE budget terms at the wake scale are plotted in figure 5.16*e-h*. The amplitudes of the TKE budget terms are approximately one-tenth of those at the *monami* scale (figure 5.16*a-d*). The waving term as a source term is as significant as the shear production term, the dominant sink is the viscous dissipation term, and the inter-scale transport is a sink inside the canopy and a source above the canopy.

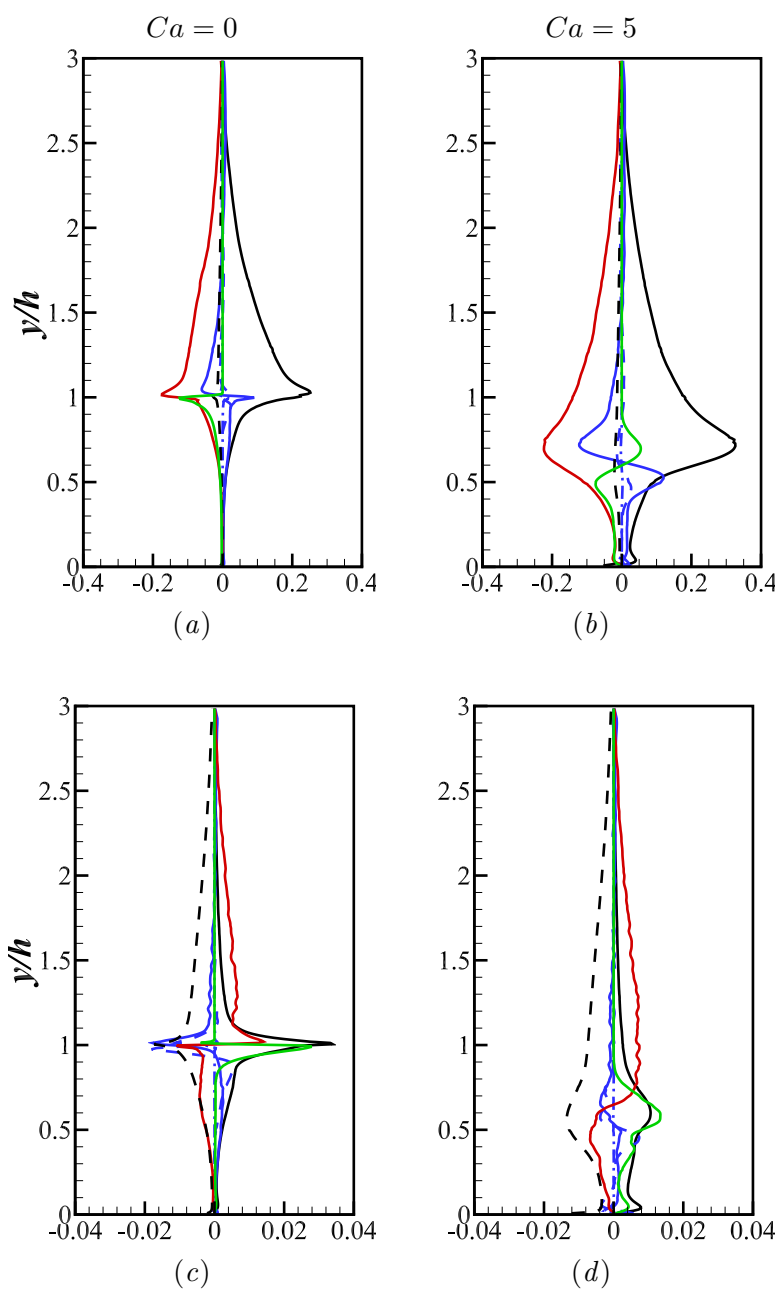
The DKE budget can be calculated through the difference between the budgets of the following two TKE definitions:

$$e_{TKE,1}(x, y, z, t) = \frac{1}{2} (u_i - \bar{u}_i) (u_i - \bar{u}_i), \quad (5.23a)$$

$$e_{TKE,2}(x, y, z, t) = \frac{1}{2} (u_i - \langle \bar{u}_i \rangle) (u_i - \langle \bar{u}_i \rangle), \quad (5.23b)$$

$$e_{DKE,p}(x, y, z) = e_{TKE,2}(x, y, z, t) - e_{TKE,1}(x, y, z, t) = \frac{1}{2} (\bar{u}_i - \langle \bar{u}_i \rangle) (\bar{u}_i - \langle \bar{u}_i \rangle), \quad (5.23c)$$

where $e_{TKE,1}$ and $e_{DKE,p}$ are equivalent to e_{TKE} and e_{DKE} , respectively, without temporal and spatial averaging. Note that because DKE is time-independent by definition and that the time-averaged flow has only the wake scale and its integer multipliers, the spectral budget of DKE has only these scales. Next, we calculate the DKE budget by



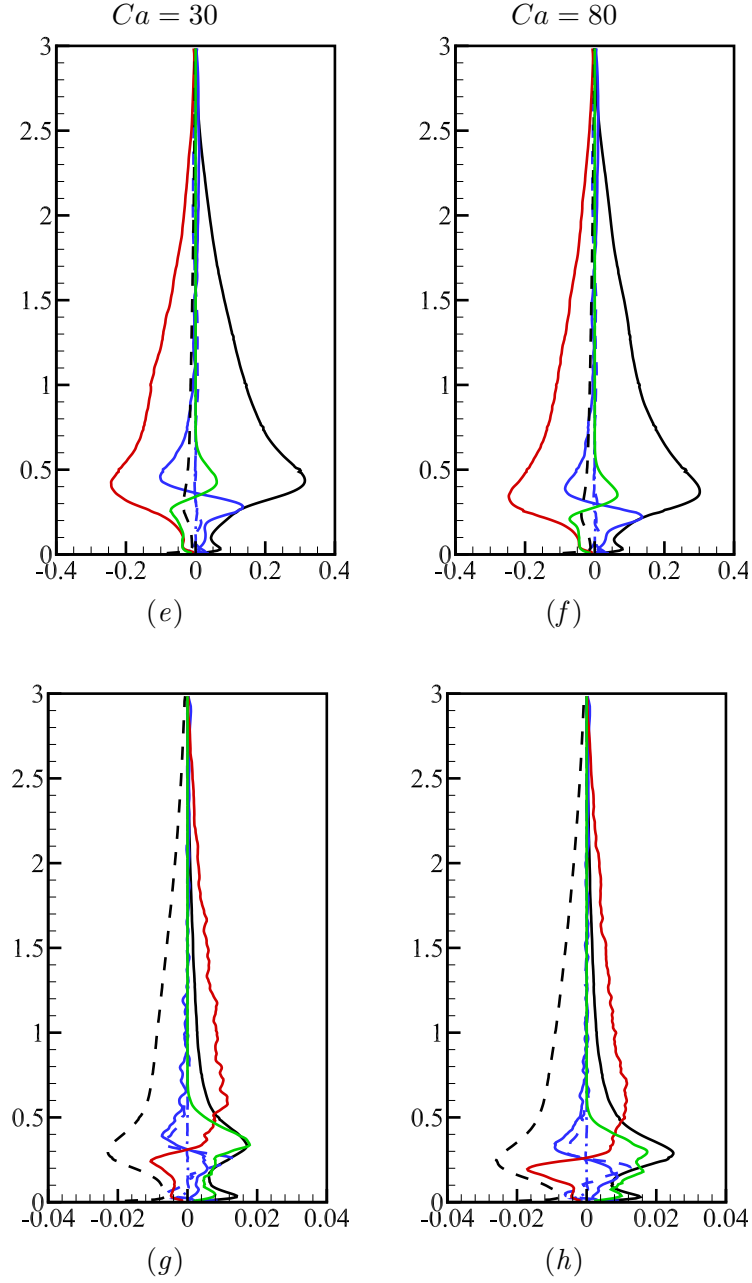


Figure 5.16: Vertical profiles of the TKE budget terms at (a,b,e,f) the *monami* scale ($k_x h \sim 0.6$) and (c,d,g,h) the wake scale ($k_x h = 4\pi$): —, the production term E_P ; —, the wall-normal turbulent transport term E_{NTT} ; - - -, the wall-normal pressure transport term E_{NPT} ; — · —, the wall-normal viscous transport term E_{NVT} ; —, the inter-scale transport term E_{IST} ; - - -, the viscous dissipation term E_{DIS} ; and —, the waving term E_W . The cases are (a,c) $Ca = 0$, (b,d) $Ca = 5$, (e,g) $Ca = 30$, and (f,h) : $Ca = 80$.

subtracting the spectral budget of $e_{TKE,1}$ from that of $e_{TKE,2}$, both of which are calculated following the algorithm in § 5.2. Therefore, the DKE budget has the same terms as the spectral TKE budget, as described in § 5.2. Figure 5.17 shows the vertical profiles of the DKE budget terms at the wake scale. Note that harmonics similar to those that occur in the TKE spectra described in § 5.3 are also observed at the integer multipliers of the wake scale and are not plotted here. Although e_{DKE} is much smaller than e_{TKE} , as shown in figure 5.1, the DKE is concentrated on the wake scale and its harmonics such that its budget terms on the wake scale (figure 5.17) are more pronounced than the TKE budget terms (figure 5.16). For the shear production term, in the flexible canopy cases, there is a local maximum at the canopy top and another local maximum near the channel bottom wall. This phenomenon is similar to the spectra of E_P for TKE at the wake scale discussed in § 5.3 and can also be explained by the vertical variation of shear. Additionally, this term is negative between the two maxima. A negative production region was also observed by [88] in their study of air turbulence over a slow water wave. Note that the TKE budget in [88] corresponds to the sum of the TKE and DKE budgets in our case, and our analyses of the TKE and DKE budgets show that the negative value is associated with the DKE budget. The wall-normal pressure transport term is the predominant term among the three wall-normal transport terms. For the rigid canopy case, the wall-normal pressure transport term is mainly responsible for transporting energy downward within the canopy; for the flexible canopy cases, this term transports DKE to the upper canopy from the lower canopy as well as from above

the canopy. As the canopy flexibility increases, the wall-normal turbulent transport term becomes significant inside the canopy. The inter-scale transport term is negative inside the canopy and transports the DKE to its harmonics. The waving term is an DKE source and is more prominent than the shear production term and is balanced mainly by the wall-normal pressure transport term in the lower canopy, the inter-scale transport term in the upper canopy, and the wall-normal turbulent transport and dissipation terms over the canopy. Summaries of the DKE fluxes at the wake scale for the rigid and flexible canopy cases are demonstrated in figure 5.18(*b*) and (*d*), respectively.

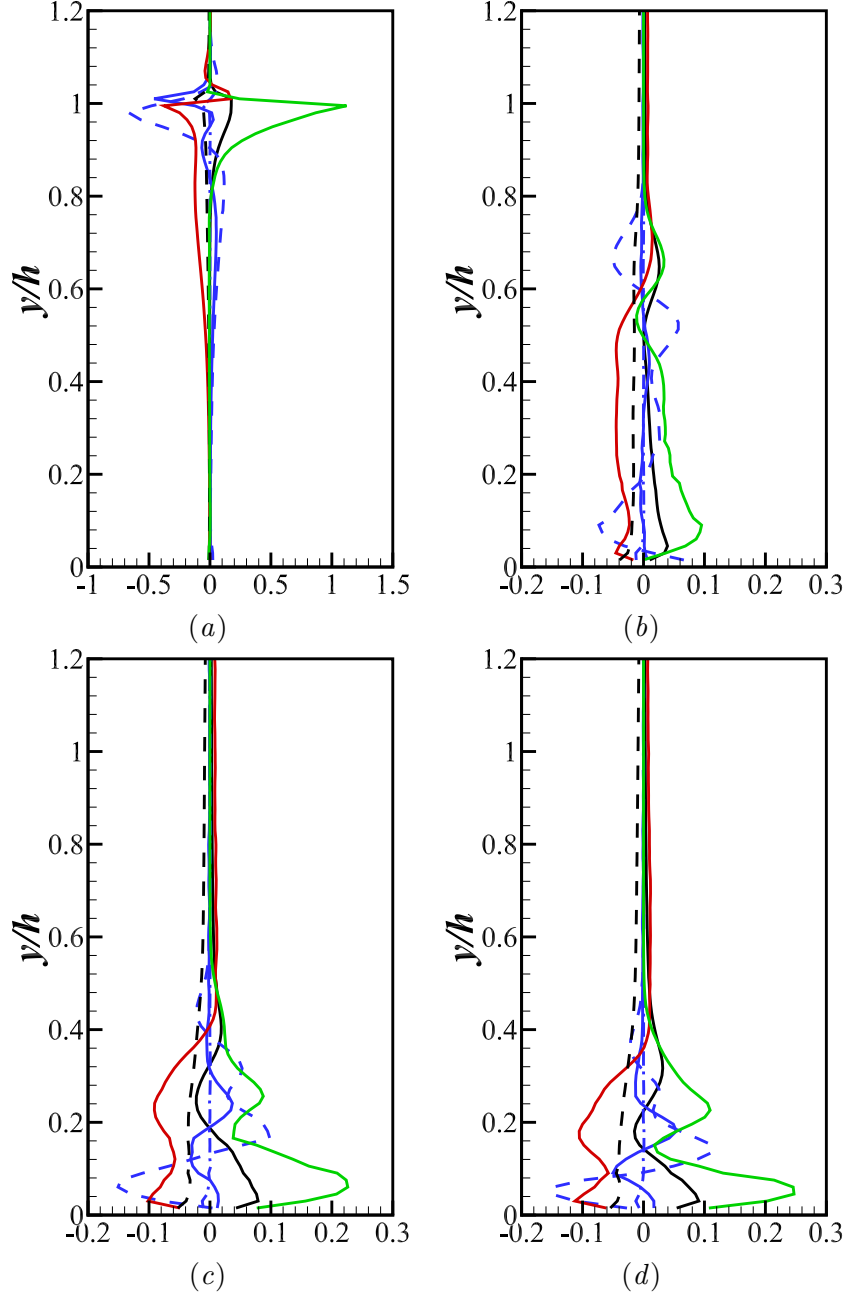


Figure 5.17: Vertical profiles of the DKE budget terms at the wake scale ($k_x h = 4\pi$) for $Ca = 0, 5, 30$, and 80 . All the terms are normalized by u^{*3}/h^3 . The lines are as follows: —, the production term E_P ; —, the wall-normal turbulent transport term E_{NTT} ; - - -, the wall-normal pressure transport term E_{NPT} ; — · —, the wall-normal viscous transport term E_{NVT} ; —, the inter-scale transport term E_{IST} ; - - -, the viscous dissipation term E_{DIS} ; and —, the waving term E_W . Note that the range of the horizontal axis differs between the case of $Ca = 0$ and the other cases.

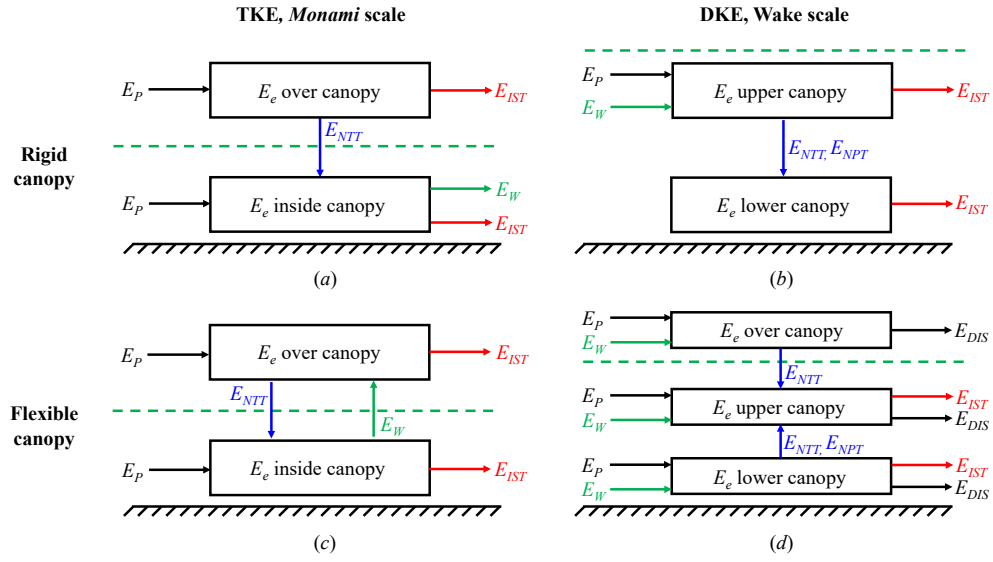


Figure 5.18: Schematics of the (a,c) TKE flow at the *monami* scale and (b,d) DKE flow at the wake scale for (a,b) rigid and (c,d) flexible canopy cases. The terms are as follows: E_P , the production term; E_{DIS} , the dissipation term; E_W , the waving term; E_{NPT} , the wall-normal pressure transport term; E_{NTT} , the wall-normal turbulent transport term; and E_{IST} , the inter-scale transport term. The horizontal green dashed line indicates the canopy top.

Chapter 6

Conclusions

This thesis contributes in two aspects. First, this thesis proposes an improvement to the conventional sharp interface IB method for simulating the interaction among air, water, and thin flexible structures. Second, this thesis simulates turbulent canopy flows by resolving individual stems with a diffused interface IB method and analyzes the flow features and the spectral TKE and DKE budgets.

6.1 Solving interaction among air, water, and flexible thin structures with sharp interface IB method

The first part of this thesis (Chapter 2) introduces an IB-based numerical algorithm for simulating the interaction between multiphase flows and thin flexible structures. The motions of two fluid phases are simulated by solving the same governing equations

with varying density and viscosity of the fluid. The free surface is captured by the CLSVOF method. The dynamics of the structure is solved by an FEM solver based on the Kirchhoff thin-shell theory. Either loose or strong coupling can be employed based on the density ratio between the fluid and structure. We propose an explicit pressure boundary condition on the immersed boundary. This is critical for obtaining the correct flow momentum flux across the immersed boundary and the correct pressure distribution around the immersed boundary. For the interaction between multiphase flows and thin flexible structures, this method ensures no fluid leakage across the thin structures. We also developed a new scheme for categorizing the Eulerian grid node to avoid extra pressure interpolation.

The proposed algorithm is tested with two 2D and one 3D canonical problems. All the cases show that explicitly applying the pressure boundary condition on the immersed boundary can improve the numerical prediction on plate deformation, and the improvement is more prominent for the 3D case. Also, in the 2D case, we demonstrate that applying the pressure boundary condition on the immersed boundary can improve the numerical result of pressure gradient around the immersed boundary.

6.2 Canopy flow simulation

In the second part of this thesis (Chapter 3–5), we simulate turbulent canopy flows with the canopy exhibiting various degrees of flexibility and investigate the flow features and the correlations between the flow and canopy dynamics. Different from the traditional

approach that models the canopy as a continuous medium with a homogeneous drag coefficient, in this study, the hydrodynamic effect of the canopy is modeled by an IB method in which individual stems within the canopy are resolved. By using this approach, *a priori* drag coefficient models are not needed, and the complex kinematics of the canopy, such as the waving motion of the stems and the phenomenon known as ‘*monami*’, are resolved. The flow is simulated by LES, and the turbulent flow motions and motions of the canopy stems are dynamically coupled with two-way interactions. The Cauchy number Ca , which denotes the ratio between the hydrodynamic force on the stem and the stem restoring force, is varied from 0 to 80, which covers a broad range of stem flexibilities from rigid stems to oscillatory stems to stems yielding to the flow.

The features and underlying mechanisms of the flows and canopies are studied. Consistent with the conclusions of previous experimental studies, our simulations reveal correlations between the turbulence structures in the mixing layer at the canopy top and the wave phase of *monami*. We find that the stem fluctuation amplitude decreases as Ca increases because highly flexible stems yield more to the incoming flow; we call this phenomenon ‘high flexibility-induced rigidity’ and use it to explain the similarities of some flow features between the cases of $Ca = 0$ and $Ca = 80$, such as the similarity among the DKE profiles. By extracting the dispersion relation of *monami* in the flexible canopy cases from the frequency–wavenumber spectrum of the canopy envelope, we find that the *monami* wave speed c is in the neighbourhood of $2u^*$. Therefore, some features of the flow over the canopy, such as the cat’s eye structure of streamlines and the

critical layer near the canopy top in the *monami*-following frame, are similar to those in the wind over a slow water wave. To understand the flow features correlated with the *monami* wave phase, we define four categories of conditionally averaged events according to the stem kinematics (figure 4.6) and perform conditional averaging on the flow velocity and drag force. We discover that the forward and backward strokes of stems correspond to sweep and ejection events, respectively. A summary of the flow pattern in the *monami*-following frame is provided in figure 4.9.

Furthermore, analyses are performed on the TKE budget. In addition to the terms describing pure channel flows, three new terms are introduced to the TKE budget equation owing to the canopy: the wake production term and dispersive transport term are related to the spatial inhomogeneity, while the waving term is related to the correlation between the hydrodynamic drag and the stem waving motion. The shear production term is the dominant energy source, but the wake production term is more pronounced inside the canopy. The turbulent transport term is the strongest among all four transport terms and plays an important role in balancing the high energy production above the canopy with the high dissipation inside the canopy. While the role of the waving term was not clarified in previous studies, our study shows that the waving term is a major TKE source term and that its amplitude can be as large as one-half of that of the shear production term in flexible canopy cases. In addition to quantifying TKE, we also quantify DKE and report that the amplitude of DKE is approximately 10% of that of TKE, which is consistent with the findings of previous

experimental studies.

In addition, for the first time, this study conducts spectral TKE budget analyses for turbulent canopy flows. Two streamwise scales are identified, namely, the wake scale associated with the interval between adjacent stems and the *monami* scale associated with the coherent structures in the mixing layer. At the wake scale, compared to the rigid canopy case, the flexible canopy cases have an additional region characterized by prominent shear production near the bottom, which is due to the strong shear at the bottom in the flexible canopy cases. The wall-normal pressure transport term has a larger amplitude than the wall-normal viscous transport term, which contrasts the case of wind over a slow water wave studied by [88]. This discrepancy is due to the different relative amplitudes of the pressure drag and viscous drag terms. Over the canopy, the inter-component transport term necessarily transports TKE from the streamwise component to the other two components, while at scales larger than the wake scale, TKE flows from the spanwise component to the vertical component. Inside the canopy, there is a weak TKE flow from the vertical component to the other two components. Additionally, the inter-scale transport term extracts the TKE at the wake and *monami* scales and redistributes it to smaller scales following a normal energy cascade. Therefore, the inter-scale transport term is unlikely to be responsible for the spectral shortcut mechanism. Moreover, a triadic analysis indicates that no inverse energy cascade occurs in our simulation cases.

This study is able to quantify the dynamic role of the energetic waving terms by

leveraging the present stem-resolving computational approach. The waving terms transfer MKE to DKE and TKE. A summary of the energy flow among MKE, DKE, and TKE and the roles of the corresponding waving terms is illustrated in figure 6.1. The streamwise component is the primary contributor to the waving term. In addition, the vertical component contributes to the DKE waving term near the canopy top in the case of $Ca = 5$. The waving term displays the effect of an inter-scale transport term in the rigid canopy case and the combined effects of the inter-scale transport term and wall-normal transport term in the flexible canopy cases. The waving term transports energy from the *monami* scale to the wake scale, indicating that this term is responsible for the spectral shortcut mechanism. Accordingly, quadrant analyses are conducted on the correlation between the drag and velocity fluctuations to understand the contributions of various *monami* wave phases to the waving term. For the correlation in the streamwise direction, $f'_1 u'_1$, large deformation and large forward speed events exhibit positive contributions owing to Q1 events, while large backward speed events yield positive contributions owing to Q3 events, and small deformation events provide positive contributions near the stem tips and negative contributions in the upper canopy owing to Q3 and Q4 events, respectively. For the correlation in the vertical direction, $f'_2 u'_2$, large deformation and large forward speed events have positive and negative contributions near the canopy top owing to Q3 and Q4 events, respectively, whereas small deformation and large backward speed events exhibit positive and negative contributions in the upper canopy owing to Q1 and Q2 events, respectively. We also calculate

the DKE budget at the wake scale. Summaries of the TKE flow at the *monami* scale and the DKE flow at the wake scale are shown in figure 5.18.

In conclusion, by resolving the dynamics of individual stems with an IB method, the present work investigates the dynamics of flows and canopies with various stem flexibilities. Our analysis of the spectral TKE budget identifies the *monami* scale and wake scale as the characteristic scales and reveals the significant role of the waving term in the spectral shortcut mechanism.

6.3 Future studies

6.3.1 Incorporating contact line model into IB method for thickness-unresolvable structures

In our algorithm for the interaction among air, water, and flexible thin structures, we assume high Weber and Capillary numbers, such that the interfacial surface tension can be neglected. When Weber/Capillary number is low, surface tension can significantly influence the flow regime and should be included [123]. In the cases of free surface intersecting with structure surface, the appearance of air–water–structure interface, also known as the contact line, imposes great challenge to the modeling [124, 125]. In recent years, a handful of algorithms incorporating contact line models into IB methods have been developed to take advantage of the IB method in modeling structures with complex geometry and arbitrary relative motion [126, 127, 128, 129, 130, 131, 132, 133].

However, these algorithms require extrapolation the level-set or volume fraction function into structures to set the contact angle on the contact line, which implies that these methods cannot be directly applied to our algorithm for thickness-unresolvable structures. Therefore, we believe that the incorporation of the contact line model into our IB method for thickness-unresolvable structures should be conducted in future studies.

6.3.2 Multiphysics processes in canopy flow

The present simulations of canopy flow are adiabatic and without particle and scalar transport. While in many industrial and natural cases, heat transfer and scalar/particle transport are significant physical processes. For example, the bathymetry-dependent sun heating and evaporation cooling over coral reefs can lead to baroclinic pressure gradient and induce exchange flows across the coral reefs [134]. The thermal status of atmospheric boundary layer and human activities in heat production strongly influence the airflow patterns in urban areas [67]. In terms of scalar/particle transport, it is found that the canopy density affects the concentration of fine particles and organic content in the underlying sediment [63], and in the natural canopy flow the dispersion of fine particles (fungal spores, pollens, small seeds, etc.) and scalars (water vapour, ammonia, ozone, etc.) is strongly driven by the turbulent flow inside the canopy [e.g. 135, 136, 137]. In view of this, we believe that the multiphysics processes should be introduced in our stem-resolved canopy flow simulation and explored in future studies.

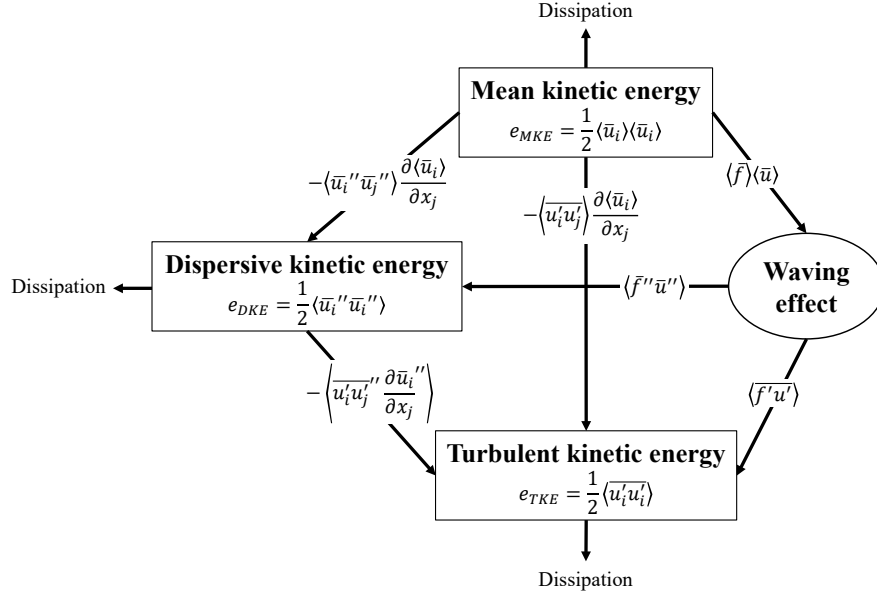


Figure 6.1: Schematics of the energy flow among MKE, DKE, and TKE and the role of the waving term.

6.3.3 Parameterization models for large-scale canopy flow simulation

The present large-scale atmospheric and ocean models [e.g. 138, 139] usually model the underlying plant canopy as surface roughness and apply turbulent boundary layer model such as Monin–Obukhov similarity theory [140]. However, it is known that the transport of momentum, heat, scalar and particle transport across the canopy top is dominant by the coherent structures at the mixing layer, not the turbulent boundary layer. Therefore, parameterization models for these transport processes for large-scale models where the canopy cannot be resolved should be developed in future studies.

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