George K. El Khoury

Numerical simulations of massively separated turbulent flows
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Norwegian University of Science and Technology
Faculty of Engineering Science and Technology
Department of Marine Technology
Numerical simulations of massively separated turbulent flows.

George K. El Khoury 2010
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Abstract

It is well known that most fluid flows observed in nature or encountered in engineering applications are turbulent and involve separation. Fluid flows in turbines, diffusers and channels with sudden expansions are among the widely observed areas where separation substantially alters the flow field and gives rise to complex flow dynamics. Such types of flows are referred to as internal flows since they are confined within solid surfaces and predominantly involve the generation or utilization of mechanical power. However, there is also a vast variety of engineering applications where the fluid flows past solid structures, such as the flow of air around an airplane or that of water around a submarine. These are called external flows and in the former case the downstream evolution of the flow field is crucially influenced by separation.

The present doctoral thesis addresses both internal and external separated flows by means of direct numerical simulations of the incompressible Navier-Stokes equations. For internal flows, the wall-driven flow in a one-sided expansion channel and the pressure-driven flow in a plane channel with a single thin-plate obstruction have been studied in the fully developed turbulent state. Since such geometrical configurations involve spatially developing turbulent flows, proper inflow conditions are to be employed in order to provide a realistic fully turbulent flow at the input. For this purpose, a newly developed technique has been used in order to mimic an infinitely long channel section upstream of the expansion and the obstruction, respectively. With this approach, we are able to gather accurate mean flow and turbulence statistics throughout each flow domain and to explore in detail the instantaneous flow topology in the separated shear layers, recirculation regions as well as the recovery zones.

For external flows, on the other hand, the flow past a prolate spheroid has been studied. Here, a wide range of Reynolds numbers is taken into consideration. Based on the characteristics of the various flow regimes as well as the recovery zones.

Descriptors: Direct numerical simulations, incompressible flows, separation, turbulent inflow conditions, wall-bounded flows, sudden expansion flows, obstructed flows, bluff body flows, prolate spheroid.

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This thesis is divided into two main parts. The first part consists of an introduction and a brief review of the main results obtained. The second part includes the papers and is divided into two sections: wall-bounded flows and bluff body flows.

The work on wall-bounded flows was in collaboration with the Department of Energy and Process Engineering at NTNU where MB was a PhD student at that time. BP and HIA supervised the work and assisted in the analysis and interpretation of the findings and in the writing process.

The thesis is based on and contains the following papers.


List of Publications

Journal Papers

Meth. Fluids 60, 263-274.

Paper 2. Barri, M., El Khoury, G. K., Andersson, H. I. & Pet-
tersen, B. 2009 Inflow conditions for inhomogeneous turbulent flows. Int. J.

Paper 3. Barri, M., El Khoury, G. K., Andersson, H. I. & Pet-
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Crossflow past a prolate spheroid at Reynolds number of 10 006. J. Fluid

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Fluid flows comprise a wide variety of complex flow phenomena which depend on the geometrical configuration of the problem as well as on the Reynolds number of the flow ($Re$). Based on the value of the latter non-dimensional parameter, a fluid motion can be classified as laminar or turbulent. Laminar flows exist at relatively low Reynolds numbers and, in this case, the motion of the fluid particles can be predictable and along observable path. However, it is well known that most fluid flows observed in nature or encountered in engineering applications are turbulent. This is because many gases and liquids have extremely low viscosity and, therefore, most practical flows are characterized by large values of the Reynolds number.

There is merely no formal definition of turbulence. According to many researchers, it is simply better to note that when the inertial forces are large in comparison to viscous forces (low viscosity), the flow develops a chaotic, random motion and is highly unpredictable. In this case, the flow field is very sensitive to perturbations and fluctuates wildly in time and space. Additionally, it contains swirling flow structures (eddies) with characteristic length, velocity and time scales which are spread over broad ranges. The size of the largest eddies in a turbulent flow is determined by the characteristic length scale of the mean flow whereas for the smallest eddies, their size depends on the Reynolds number. At high Reynolds numbers, there is a wide range of eddy sizes. The largest eddies, which are generated directly by shear in the mean flow, contain most of the turbulent kinetic energy and are unaffected by viscous stresses. However, due to inertial effects these large eddies break-down and transfer energy to smaller eddies that by its turn spread the energy to yet smaller structures. This process continues until a scale is reached at which the inertial forces become comparable to the viscous stresses and the kinetic energy is dispersed by the latter into heat.

### 1.1. Navier-Stokes equations

The spatial and temporal evolution of a viscous fluid can be described by the Navier-Stokes equations. These equations are derived by applying Newton's second law to a continuum, the principle of mass conservation and a constitutive law which relates the shear stresses in a fluid to the rate of deformation of a fluid element. In the present thesis, we consider a linear stress-strain relation and thus the fluid is Newtonian. Based on these assumptions and by introducing heat.

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considering an incompressible isothermal fluid, the Navier-Stokes equations are expressed as:

\[
\frac{\partial u_i}{\partial t} + \frac{\partial }{\partial x_j} \left( \rho u_i u_j \right) = - \frac{\partial p}{\partial x_i} + \nu \frac{\partial^2 u_i}{\partial x_j \partial x_j} + \frac{\partial \tau_{ij}}{\partial x_j}, \quad i, j = 1, 2, 3 \tag{1.1}
\]

Here, \( \rho \) and \( \nu \) are the density and the kinematic viscosity of the fluid, respectively. It is often convenient to work with a non-dimensional form of the Navier-Stokes equations by introducing representative scales for velocity, length, time, and pressure, \( U, L, T, \) and \( p^* \) (the velocity and length scales have to be chosen appropriately to the flow problem considered), a non-dimensional form of the Navier-Stokes equations is obtained and a Reynolds number can be defined as:

\[
Re = \frac{UL}{\nu} \tag{1.3}
\]

The Reynolds number can be interpreted as a measure of the inertial forces divided by the viscous forces, and is by far one of the most important non-dimensional numbers in fluid mechanics.

### 1.2. Direct numerical simulation

Direct numerical simulation (DNS) is perhaps the most straightforward approach available to numerically solve the Navier-Stokes equations. It involves no turbulence modelling and directly solves the equations of fluid motion capturing all eddies in the flow field ranging from the characteristic length scale (\( L \)) right down to the Kolmogorov length scale (\( \eta \)) related to the smallest eddies. The Kolmogorov scales represent the smallest turbulent motions and are defined as:

\[
\eta \equiv \left( \frac{\nu}{\epsilon} \right)^{1/4} \tag{1.4}
\]

\[
\tau_{ij} \equiv \left( \frac{\nu \epsilon}{\lambda} \right) \tag{1.5}
\]

\[
u = \left( \frac{\nu \epsilon}{\lambda} \right) \tag{1.6}
\]

for length, time, and velocity, respectively. Here, \( \epsilon \) denotes the turbulent kinetic energy dissipation. By considering an inviscid estimate for the dissipation rate, i.e., \( \epsilon \sim U^3/L \) and the Reynolds number introduced in the previous subsection, the following ratio between the smallest and largest scales is obtained:

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### 1.3. Introduction

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\frac{\eta}{L} \sim \frac{1}{Re} \tag{1.7}
\]
It is apparent from the above relation that the separation in scales increases dramatically with the increase in the Reynolds number. In fact, the number of grid points needed to perform a three-dimensional DNS is proportional to $Re^{9/4}$. By taking into consideration the total increase in the number of time steps associated with the increase in $Re$ ($L/U\tau_\eta \sim Re^{1/2}$), it can be concluded that the computational cost in DNS roughly grows by $Re^{3}$. With such high computational costs, DNS is mainly used as a research tool to study moderately low Reynolds number flows.
Inflow conditions

The generation of realistic dynamic inflow conditions for spatially developing turbulent flows is among the most fundamental issues in direct numerical simulations. Accordingly, it is hardly surprising that this problem has attracted considerable interest among researchers since the usage of periodic boundary conditions is severely restricted or not suitable in many practical geometrical configurations. This is because of the flow inhomogeneity in more than one direction. Several methods of different complexity have been previously investigated in order to provide realistic inflow data for inhomogeneous turbulent flows.

In numerical simulations of turbulent boundary layers, Spalart & Watmuff (1990) solved the incompressible Navier-Stokes equations over a flat plate using periodic boundary conditions in the x and z directions parallel to the plate. Although such a flow problem in not homogeneous in the z-direction, the authors made it feasible to use periodic boundary conditions in the latter direction by introducing a fringe region technique. In this method, a forcing term is applied on the Navier-Stokes equations and within a certain period in the they distinguished between a fringe region, close to the inflow and outflow boundaries, in which the added terms are finite and a useful region in which they are zero. Subsequently, the flow profile close to the outflow is returned to the desired inflow profile in the fringe region. Perhaps one of the most studied flow configurations in which the turbulence then develops. However, this requires computational domain in order that realistic turbulence develops properly from the added perturbations. Nonetheless, other methods have been employed by researchers, for instance, proper orthogonal decomposition (POD).

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A more realistic approach to generate proper turbulent inflow conditions for spatially developing channel flows would involve the utilization of an auxiliary simulation. In this case, the auxiliary simulation is run with streamwise periodicity and synchronously with the actual computation under consideration. Thereafter, instantaneous velocity fields are taken from the auxiliary simulation and simultaneously fed at an inflow of the main simulation in order to mimic a long channel section upstream the inlet of the spatially developing problem. However, since this method demands the run of two separate simulations simultaneously, it can be very expensive in terms of CPU time. To address this issue, a newly developed technique is introduced to generate dynamic inflow conditions for inhomogeneous turbulent flows; see Buzzi et al. (2000). The method, referred to as the cost-effective method, consists of recycling finite-length time series (ts) of instantaneous velocity profiles at the input of the simulation under consideration. These profiles are taken from a precursor simulation and a physical constraint is introduced on the length of the time series, to be of order of the large-eddy turnover-time. Figure 2.1 displays the turbulence intensities and shear stress variations obtained via the cost-effective method in the inlet section of DNS studies of turbulent channel flow with an obstruction and turbulent Couette flow over a backward-facing step. The agreement between the simulation and the DNS data of Kim et al. (1987) is rather satisfactory which demonstrate the ability of the present technique to generate suitable inflow conditions for inhomogeneous turbulent flows.

Figure 2.1. Cost-effective method applied to inhomogeneous wall-bounded flows. (a) Comparison with DNS data from Kim et al. (1987) for fully developed turbulent Poiseuille flow. Turbulent intensities: – – – –, streamwise direction; ······, wall-normal direction; ⋄, spanwise direction. (b) Comparison with DNS data from Bech et al. (1995) for fully developed turbulent Couette flow. Shear stress variations: – – – –, viscous shear stress; – – – –, turbulent shear stress; ······, total shear stress ϱ.

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3.1. One-sided expansion channels

Due to its geometrical simplicity, the flow over a backward facing step (BFS) has been extensively investigated through the past few decades in order to study wall-bounded separated flows. Typical prototypes of BFS flows are the boundary layer, the plane channel and the Couette flow cases; see e.g. Eaton & Johnston (1981) and Le et al. (1995). A common feature of these flows is the existence of a shear layer emanating from the step corner and reattaching further downstream leading to the formation of a recirculation bubble. The presence of the internal shear layer and the massive recirculation zone gives rise to complex flow dynamics which for instance affect the turbulence production and Reynolds stress anisotropy.

Nevertheless, the most studied BFS flow is the pressure-driven flow in a plane channel with a sudden one-sided expansion; see e.g. Barrit et al. (2010). Owing to the principle of mass conservation, the Reynolds number remains the same downstream of the step as in the upstream part of the channel. In a BFS Couette flow, on the other hand, the Reynolds number becomes higher downstream of the step. It is well known that the shear-driven turbulent Couette flows (Bech et al. 1995) exhibit a number of characteristic features which make them distinguishingly different from the pressure-driven Poiseuille flow, notably the monotonically increasing mean velocity profile.

![Figure 3.1. Instantaneous enstrophy contours revealing the formation of Kelvin-Helmholtz instabilities in turbulent Couette backward-facing step flow.](image)

Wall-bounded flows

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![Figure 3.1. Instantaneous enstrophy contours revealing the formation of Kelvin-Helmholtz instabilities in turbulent Couette backward-facing step flow.](image)
Since the turbulent flow field in BFS Couette flow evolves spatially in the streamwise direction, special attention must be paid to the boundary condition at the input. For this purpose, the cost-effective method introduced in the previous chapter is utilized to generate realistic dynamic inflow conditions from a precursor simulation. Depicted in figure 3.2 is the mean velocity profile at two representative locations: in the upstream channel close to the inlet, and in the recovery region far downstream the step. At the former position, excellent agreement is observed with the DNS data from Bech et al. (1995), while at the latter position the profile is yet far from being anti-symmetric. Midway between the walls $U_z$ is still roughly half of $U_{1/2}$ which should have been reached in the case of a fully redeveloped Couette flow. Nonetheless, irrespective of the length of the domain that can be used in the downstream part of the channel, an anti-symmetric profile corresponding to a fully redeveloped Couette flow will not be reached. This is due to the principle of mass conservation. It follows that since the height of the domain after the step is twice that of the inlet section and the mean velocity profile of Couette flow is monotonically increasing to a constant value of $U_w$, the flow cannot adjust itself to an anti-symmetric $S$-profile shape and at the same time maintain a constant flow rate.

This observed pattern is a distinctive feature in the BFS Couette flow. The two-dimensional upstream mean flow develops to an essentially uni-directional flow in the downstream part of the computational domain, i.e. beyond $y/h \approx 30$ or 15 step heights $h$ downstream of the sudden expansion. It is noteworthy that the upstream pure Couette flow redevelops into a mixed Couette-Poiseuille flow in contrast to the classical pressure-driven backward-facing step flow where an upstream Poiseuille flow inevitably redevelops to another pure Poiseuille flow far downstream of the step. In the present case, however, an adverse pressure gradient is established with the view to assure global mass conservation. The resulting mixed Couette-Poiseuille flow cannot adjust itself to an anti-symmetric $S$-profile shape and at the same time maintain a constant flow rate.

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flow exhibits major asymmetries in the turbulence field with a substantially reduced turbulence level along the stationary wall, and closely resembles the Couette-Poiseuille flow simulations reported by Kuroda et al. (1995).

In order to study the shear layer which forms when the upstream Couette flow mixes with the recirculating flow downstream of the step, the turbulent kinetic energy budget is presented at a streamwise position passing through the center of the primary separation zone (i.e. 3h downstream of the step).

The transport equations for the Reynolds stress tensor are:

\[
\frac{D}{Dt} \tau_{ij} = P_{ij} - \varepsilon_{ij} + \Pi_{ij} + G_{ij} + D_{ij} + T_{ij}
\]  (3.1)

where the production, dissipation, pressure-strain, pressure diffusion, molecular diffusion and turbulent diffusion are defined as:

\[
P_{ij} = -\frac{\partial u_i}{\partial x_j} - \frac{\partial u_j}{\partial x_i}
\]

\[
\varepsilon_{ij} = 2\tau_{ij} \frac{\partial u_i}{\partial x_j}
\]

\[
\Pi_{ij} = \frac{1}{\rho} \left( \frac{\partial u_i}{\partial x_j} - \frac{\partial u_j}{\partial x_i} \right)
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The budget for the turbulent kinetic energy $\frac{\partial^2}{\partial x^2} = \frac{\partial^2}{\partial y^2}$ is one half the sum of the budget of the diagonal components of the Reynolds stress tensor, and is shown in figure 3.3. In the shear layer, production and dissipation are the most dominant terms whereas turbulent diffusion transports energy into the upper half part of the channel only. The contribution to $P_{ij}$ comes mainly from $P_11$ (see figure 3.4) and its peak at $y/h \approx 1$ is almost 2.5 times larger than that of $\epsilon$. This infers that dissipation is not in balance with production in the regions close to the wall, turbulence is substantially damped along the lower one due to recirculation. In this case, the production and turbulent diffusion terms are almost negligible which is in contrast to viscous dissipation that is the most significant (among the plotted terms).

Since there is nearly no production of $\epsilon$, as can be inferred from figure 3.4, their only source of energy is from $\Pi_{11}$ which serves to redistribute energy between the normal stresses. This is shown in figure 3.5 where the pressure-strain terms appearing in the $\pi_{11}$, $\pi_{22}$ and $\pi_{22}$ equations are plotted together at $x/h = 18$. Across the channel, $\pi_{11}$ sets as a receiving component taking energy mainly from $\epsilon$. The profiles of $\Pi_{11}$ and $\Pi_{22}$ are plotted together at $x/h = 18$ and its peak at $y/h \approx 1$ is almost 2.5 times larger than that of $\epsilon$. This infers that dissipation is not in balance with production in the regions close to the wall, turbulence is substantially damped along the lower one due to recirculation. In this case, the production and turbulent diffusion terms are almost negligible which is in contrast to viscous dissipation that is the most significant (among the plotted terms).

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indicate a qualitative difference in the energy exchange pattern between the two walls. While there is a large energy transfer from the $i$ component to the $j$ and $k$ component between $y/h = 0$ and $y/h \approx 0.2$, $i$ contributes to $j$ in delivering energy to $k$ in the region near to the upper wall. Away from the walls, the major effect of the pressure-strain is to distribute energy from $i$ component to the other two components. The sum of the three components (i.e. $\Pi_i$) is almost zero and this supports the adequacy of the sampling procedure.

### 3.2. Plane channel with an obstruction

Asymmetric flow patterns that develop in symmetric channel configurations are of considerable interest in a variety of industrial applications. Turbulent mixing and the rate of heat transfer are for instance among the physical mechanisms affected by flow asymmetry. Perhaps the most studied geometrical configuration related to such flows is the flow in a plane channel with a double-sided symmetric expansion; see e.g. Durst et al. (1974). A somewhat related problem in a plane channel configuration with a single or a sequence of thin-plate obstructions. In a similar approach to the previous study, the cost-effective method is used to provide realistic fully turbulent inflow conditions in order to study the asymmetric flow pattern caused by a single thin-plate obstruction in a plane channel. In this case, however, the precursor simulation is that of a fully-developed turbulent channel flow. The resulting flow field upstream the obstruction compares almost perfectly well with data from Kim et al. (1987). It is worthwhile to note that the

![Figure 3.5. Pressure-strain terms normalized by $u_i^2/\nu$ at $x/h = 18$.](image)

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current DNS study focuses on the asymmetric flow pattern in the fully developed turbulent state and not on the source of or mechanism for creating the asymmetry phenomenon.

In the present findings, the flow downstream the obstruction is asymmetric with an extraordinarily long separation bubble at one of the walls. This separation bubble extends about 17 obstruction heights downstream of the thin-plate, and is almost 4 times longer than the shorter bubble found on the other wall. After the distinctly asymmetric pattern has been established as in figure 3.6(a), the jet remains deflected to the same side due to the substantially lower pressure within the shorter recirculating eddy. A previous DNS study on a similar problem was carried out by Makino et al. (2008) but with a periodically repeating obstruction at a slightly lower Reynolds number. For that purpose, the authors used periodic boundary conditions in the streamwise direction. In comparison with their findings, the longest of the two bubbles found in the present case exceeds the downstream length of the two bubbles found in the present case exceeds the downstream length in the streamwise direction. In comparison with their findings, the longest of the two bubbles found in the present case exceeds the downstream length.

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of the domain used by Makino et al. (2008). This verifies the necessity of proper inflow conditions if the effects of a single obstruction is to be studied.

An insight on the complex motion in the present study is obtained from figure 3.6(b) where stream ribbons of the instantaneous velocity components are shown in an upstream perspective view from $x/h = 20$. Stream ribbons illustrate the direction of the flow and are constructed as narrow surfaces defined between two adjacent streamlines. Large-scale motions are observed in the lower-half of the channel which is in contrast to the upper-half where the flow is relatively calm. These swirling motions transport fluid between the primary recirculation region on the lower wall and the high-speed jet and simultaneously influence the Reynolds stresses. Meanwhile, the formation of Kelvin-Helmholtz vortices is visible from figure 3.7. They are generated around the top edges of the slits and swept away forming a fine-scale turbulence in the mixing layers. The absence of vorticity between the two mixing layers indicates a potential-like type flow (i.e. irrotational flow) in this region which is due to flow acceleration.

In a quick look at the momentum balance, the streamwise momentum equation for a channel flow with single thin-plate obstruction can be written as:

$$\rho \frac{DU}{Dt} = \frac{1}{2} \frac{DP}{\rho} \frac{\partial U}{\partial x} \frac{\partial U}{\partial y} \frac{\partial U}{\partial y} + \frac{\nu}{2} \left( \frac{\partial^2 U}{\partial x^2} + \frac{\partial^2 U}{\partial y^2} \right) \tag{3.8}$$

Here, the left hand side of equation 3.8 denotes the convective accelerations, and the term with $\partial P/\rho$ is ignored since the flow is considered to be statistically steady. The above equation can be rearranged as follows:

$$0 = -\frac{1}{\rho} \frac{\partial P}{\partial x} \frac{\partial U}{\partial x} \frac{\partial U}{\partial y} \frac{\partial U}{\partial y} \frac{\nu}{2} \left( \frac{\partial^2 U}{\partial x^2} + \frac{\partial^2 U}{\partial y^2} \right) \tag{3.9}$$

In the present case, $-\nu U$ is a result of the effect of the slits on the mean flow field as shown in figure 3.8(a). As the flow contracts, $U$ is positive whereas $V$ attains positive and negative values in the lower and upper parts of the domain used by Makino et al. (2008). This verifies the necessity of proper inflow conditions if the effects of a single obstruction is to be studied.

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of the channel, respectively. Downstream the constriction, the sign of $V$ depends on the direction in which the flow bends. Since the flow bends upwards in the present case, the positive values of $U$ and $V$ give a negative $-UV$. In figure 3.8(b), the pressure gradient $\partial P/\partial y$ attains substantial negative values in the mixing layer downstream of the upper slit, i.e. $P$ is lower in the recirculation zone than in the jet flow. The upward pressure force $-\partial P/\partial y$ therefore tends to bend the mixing layer towards the upper wall.

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{figure3.8.png}
\caption{(a) Contours of the product $-UV/U^2_i$. (b) Contours of wall-normal pressure gradient $\partial P/\partial y$.}
\end{figure}
4.1. Flow past a prolate spheroid

A wake develops behind a bluff body if the Reynolds number exceeds a certain critical value such that the flow separates from the surface of the body. The resulting wake flow may be steady or unsteady and even become turbulent depending on the actual value of the Reynolds number $Re$. The shape of the body itself, as well as its orientation relative to the oncoming flow, determines the symmetry of the resulting wake. The two most frequently studied bluff bodies are the sphere and the circular cylinder (see e.g. Johnson & Patel 1999 and Zdravkovich 1997) which at fairly low Reynolds numbers give rise to axisymmetric and planar wakes, respectively.

A prolate spheroid is a three-dimensional body with two different length scales, one plane of symmetry, and one axis of symmetry. The ratio between the semi-major and semi-minor axes, i.e. the aspect ratio, is a measure of departure from a spherical body. Spheroids with aspect ratios 8:1, 6:1, and

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3.1 can be considered as simplified models of submarines, unmanned underwater vehicles, missiles, airships, etc. When the aspect ratio is relatively high and the major axis is aligned with the flow the prolate spheroid behaves as a slender body. If, on the other hand, the inflow is perpendicular to the symmetry axis, i.e. parallel with the equatorial plane, the spheroid plays the role of a bluff body and the wake can be expected to possess two mutually perpendicular symmetry planes. The wake behind this spheroid shares some features of the cylinder wake and others of the wake of a sphere.

The wake behind a prolate spheroid without incidence is axisymmetric, at least at low Re, just as is the wake of a sphere. The axisymmetry is broken by the incidence angle and in the particular case of 90° angle of attack it can be anticipated that the wake at mid-span will resemble the wake behind circular cylinders. In the Reynolds number range up to 300, Shaoard et al. (2008) reported that the vortices shed in the vicinity of the cylinder mid-span resembled Kármán vortices when the length-to-diameter ratio exceeded 4. The helical-like wake which may occur behind a sphere at higher Re due to the shedding of hairpin vortices at varying azimuthal locations will probably be prohibited in a wake behind a prolate spheroid in crossflow provided that the aspect ratio is sufficiently above 1. At fairly low Reynolds numbers, a toroidal vortex ring develops on the rear side of a sphere; see for instance Achenbach (1974) and Johnson & Patel (1989). In the present case, however, the elliptical rather than circular crossflow provided that the aspect ratio is sufficiently above 1.

At fairly low Reynolds numbers, a toroidal vortex ring develops on the rear side of a sphere; see for instance Achenbach (1974) and Johnson & Patel (1989). In the present case, however, the elliptical rather than circular
cross-section of the spheroid severely breaks the axisymmetry of the vortex ring and leaves the stretched vortex structure shown in figure 4.3(c). The ‘croissant’-like shape results from the recirculating flow above the meridional plane \( y/D = 5.5 \), while a mirror-shaped vortex structure is present below the meridional plane. Due to the prevailing flow symmetries at \( Re = 50 \), no exchange of fluid between the two halves take place. The upstream view into the wake from the cross-sectional \((x, y)\)-plane, \( 4.5D \) downstream of the major axis of the spheroid in figure 4.3(b), gives another impression of the three-dimensional wake topology. The shape of the stream ribbons becomes gradually more distorted with the distance from mid-span and a stream ribbon emerging from the polar area is almost parallel with the meridional \((x, y)\)-plane.

\( 4. BLUFF BODY FLOWS \)

\( 4.3 \) Steady flow past a prolate spheroid at \( Re = 50 \) visualized by means of velocity stream ribbons. (a) Top view; (b) upstream view from \( x/D = 10 \).

\( 4.3 \) Steady flow past a prolate spheroid at \( Re = 50 \) visualized by means of velocity stream ribbons. (a) Top view; (b) upstream view from \( x/D = 10 \).

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\( 4.3 \) Steady flow past a prolate spheroid at \( Re = 50 \) visualized by means of velocity stream ribbons. (a) Top view; (b) upstream view from \( x/D = 10 \).
The complexity of the flow topology increases with Re and the asymmetry about the meridional plane is evident even for Re = 100 in figure 4.4 where a closed recirculation bubble does no longer exist as in the steady case. The upper vertical structure in figure 4.4 is considerably larger than the oppositely rotating lower vortex. The fluid which is deflected upwards as the spheroid is approach seems to follow the contours of the upper vortex. A part of this fluid is entrained into the spinning vortex whereas another part proceeds in the upstream direction and eventually feeds the lower vortex. The oncoming fluid which is deflected downwards, however, is directly fed into the counter-clockwise vortex.

The frontal area of the prolate spheroid in crossflow is elliptical. At least in the far wake, one may expect that the flow would resemble the wake formed behind sharp-edged elliptical disks. The plot of $a/U_c$ in the meridional plane in figure 4.1 gives the impression that the wake contracts to a narrow band in the spanwise direction as the flow progresses downstream. A clearer view of this phenomenon can be seen from the cross-stream slices depicted in figure 4.5. It is readily observed that the axis of the wake has switched at $x = 2D$, in the sense that the major axis of the wake now is aligned with the minor axis of the sphere for the Reynolds numbers under consideration here. Moreover, the shape of the near-wake reflected the elliptical shape of the meridional cross-section of the sphere. Some $10^5$ downstream, however, the major axis of the wake became aligned with the minor axis of the sphere. A similar axis switching was observed experimentally by Kuo & Baldwin (1967) and Kuo & Abe (1999) in the wake behind an elliptical disk. In that case, the authors reported that axis-switching occurred at about 4D downstream of the disk, i.e., significantly closer to the bluff body than in the present study. While Kuo & Baldwin (1967) avoided to present a hypothesis on the reason why the wake grew non-uniformly behind elliptical bluff...
bodies, Kiya & Abe (1999) attributed the mechanism of axis switching to different growth rates of the wake in the plane of the major and minor axis.
CHAPTER 5

Review of papers content

In this chapter, a brief summary of each of the five articles included in part two of the thesis is presented. The first three papers address wall-bounded flows while the last two are concerned with uniform flow past prolate spheroid at incidence.

Paper 1: Inflow conditions for inhomogeneous turbulent flows.
Barri, M., El Khoury, G. K., Andersson, H. I. & Pettersen, B.

In this paper, a new technique, referred to as the cost-effective method, is introduced in order to generate realistic dynamic inflow conditions for spatially developing wall-bounded turbulent flows. This method consists of recycling finite-length time series (ts) of instantaneous velocity profiles at the input of the simulation under consideration. The profiles are taken from a precursor simulation and a physical constraint is introduced on the length of the time series ts to be of order of the large-eddy-turnover-time. Excellent agreement with fully-developed channel flow statistics is observed when ts equals or exceeds the large-eddy-turnover-time scale. The present results are more realistic than those obtained with synthetic turbulence generation and at the same time substantially cheaper than running an auxiliary simulation in parallel.

Paper 2: Massive separation of turbulent Couette flow in a one-sided expansion channel.
El Khoury, G. K., Andersson, H. I., Barri, M. & Pettersen, B.
The newly developed technique presented in the preceding paper is used here in order to study the wall-driven flow over a backward-facing step by means of direct numerical simulation. This paper demonstrates that plane Couette flow at exactly the same Reynolds number as in the extensive investigation by Bech et al. (1995). An interesting feature in the present study is that the mean streamwise velocity U does not retain the characteristic S-shape typical of a pure Couette flow far downstream the step. This phenomenon is ascribed to the principle of global mass conservation, which can be fulfilled only if an adverse mean pressure gradient is established in the recovery region. This is indeed what is observed in the recovery region where a Couette-Poiseuille flow type prevails.

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Paper 2: Massive separation of turbulent Couette flow in a one-sided expansion channel.
El Khoury, G. K., Andersson, H. I., Barri, M. & Pettersen, B.
The newly developed technique presented in the preceding paper is used here in order to study the wall-driven flow over a backward-facing step by means of direct numerical simulation. This paper demonstrates that plane Couette flow at exactly the same Reynolds number as in the extensive investigation by Bech et al. (1995). An interesting feature in the present study is that the mean streamwise velocity U does not retain the characteristic S-shape typical of a pure Couette flow far downstream the step. This phenomenon is ascribed to the principle of global mass conservation, which can be fulfilled only if an adverse mean pressure gradient is established in the recovery region. This is indeed what is observed in the recovery region where a Couette-Poiseuille flow type prevails.
Paper 3: Asymmetries in an obstructed turbulent channel flow.

In this paper, the flow of a viscous fluid past a prolate spheroid has been investigated numerically at some different Reynolds numbers in the laminar flow regime. Contrary to earlier investigations, the major axis of the spheroid was oriented perpendicular to the oncoming flow. The wake flow therefore showed some similarities with the wake behind finite-length circular cylinders and yet also resembled the axi-symmetric wake behind a sphere in some other aspects. The wake behind the 6:1 spheroid remained steady at the lowest Reynolds numbers considered, \(Re = 50\) and 75. An unsteady wake was observed at \(Re = 100\), although the flow field still retained its symmetry about the equatorial plane. A symmetry-breaking occurred at a somewhat higher Reynolds number and more complex wake patterns were seen at \(Re = 200\) and 1000. Doubly inclined hairpin vortices, although from the two sides of the spheroid, were essential ingredients of the periodic flow field at \(Re \geq 100\). The contours of the nonaxisymmetric near-wake mimicked the shape of the spheroid. Following a switching of the axes of the wake some 10 minor diameters downstream, the major axis of the wake became aligned with the minor axis of the spheroid.

Paper 4: Crossflow past a prolate spheroid at Reynolds numbers of 10 000.

This paper represents a continuation of the previous work presented in paper 4 on crossflow past a 6:1 prolate spheroid. While the low Reynolds numbers effects are investigated in the previous paper, the current DNS simulation approaches to paper two, the cost-effective method is used to provide realistic fully turbulent inflow conditions from a precursor simulation. In the present study, however, the precursor simulation is that of a fully-developed turbulent channel flow. A previous DNS study on a similar problem was carried out by Makino et al. (2008) but with a periodically repeating obstruction at a slightly lower Reynolds number. For that purpose, the authors used periodic boundary conditions in the streamwise direction. In comparison with their findings, the longest of the two bubbles found in the present case exceeds the downstream length of the domain used by Makino et al. (2008). This verifies the necessity of proper inflow-outflow conditions if the effects of a single obstruction is to be studied.

Paper 5: Asymmetries in an obstructed turbulent channel flow.

This paper represents a continuation of the previous work presented in paper 4 on crossflow past a 6:1 prolate spheroid. While the low Reynolds numbers effects are investigated in the previous paper, the current DNS simulation approaches to paper two, the cost-effective method is used to provide realistic fully turbulent inflow conditions from a precursor simulation. In the present study, however, the precursor simulation is that of a fully-developed turbulent channel flow. A previous DNS study on a similar problem was carried out by Makino et al. (2008) but with a periodically repeating obstruction at a slightly lower Reynolds number. For that purpose, the authors used periodic boundary conditions in the streamwise direction. In comparison with their findings, the longest of the two bubbles found in the present case exceeds the downstream length of the domain used by Makino et al. (2008). This verifies the necessity of proper inflow-outflow conditions if the effects of a single obstruction is to be studied.
focuses solely on the subcritical Reynolds number of 10,000. At this Reynolds number, the boundary layer separated from the frontal side of the spheroid was also found to be laminar and formed an elliptical vortex sheet. In this case, however, the detached shear layer was unstable from its very inception and even the near-wake turned out to be turbulent. The Strouhal number associated with the large-scale shedding was 0.156, significantly below that of the wake of a sphere. A higher-frequency mode was associated with Kelvin-Helmholtz instabilities in the detached shear layer. The shape of the near-wake mirrored the shape of the spheroid. Some 10 minor diameters downstream, the major axis of the wake became aligned with the minor axis of the spheroid.
Acknowledgment

First of all, I would like to thank Professor Helge I. Andersson for his support, encouragement and inspiration throughout the past four years. This work would have not been possible without his continuous guidance and suggestions.

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My time in Trondheim was made enjoyable in large part due to my friends and the pleasant work environment at NTNU. For this purpose, I would like to thank my fellow colleagues and the staff at the department of marine technology for providing a pleasant and efficient work environment.

I am also grateful to my friends in Trondheim and the time spent with them.

Finally, I would like to express my deepest gratitude to my parents Kamal and Najwa, my brother Elie and his wife Abeer, and to my sister Eleine.

Part 2 Papers Part 2 Papers
Part 2.1
Wall-bounded flows

Inflow conditions for inhomogeneous turbulent flows
Barri, M., El Khoury, G. K., Andersson, H. I. & Pettersen, B.

Paper 1
Is not included due to copyright
Massive separation of turbulent Couette flow in a one-sided expansion channel

El Khoury, G. K., Andersson, H. I., Barri, M. & Pettersen, B.


Included in a special issue of selected papers from *The 6th International Symposium on Turbulence and Shear Flow Phenomena (TSFP-6).* Seoul-Korea, June 22-24, 2009.
Turbulent flow over a backward-facing step (BFS) is a simplified model of industrial applications. Although its geometry is simple, the flow physics is rich compared to other shear flows such as channel, plane Couette and the Couette flow cases, because of the existence of a shear-layer emanating from the step corner and resulting further downstream laminar to turbulent transition. The understanding of the flow over a backward-facing step was initially acquired by experiments and two-dimensional numerical simulations. The early studies were performed by Abbot and Kline (1980), Friedrich and Arnal (1990), Kuehn (1980), Durst and Tropea (1981), Ötügen (1991), and Ra (1993) who performed a numerical investigation of the coherent structures from a jet impinging on a backward-facing step with an expansion ratio of 2.1 and provided information on the relation between the Reynolds numbers and the recirculation length. The authors covered a wide range of Reynolds numbers from about 50 to 6000 and found that the flow can be divided into three regimes: a separation bubble regime, a transitional regime, and a fully turbulent regime. The findings of Kuehn (1980), Durst and Tropea (1981), Ötügen (1991), and Ra (1993) were a mean velocity profile for a flat plate turbulent boundary at the inlet, and results were obtained for Reynolds numbers from 50 to 6000. The authors also characterized the flow using Reynolds number and the reattachment length. The massive recirculation zone gives rise to complex flow separations. This technique was also used by neto et al. (1997) who provided an extensive DNS study of turbulent boundary layer backward-facing step flow using the large-eddy simulation (LES) technique. This technique was also used by Neto et al. (1997) who provided an extensive DNS study of turbulent boundary layer backward-facing step flow using the large-eddy simulation (LES) technique.

1. Introduction

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earlier study on numerical simulation of plane channel flow (see Le et al. (2007) and 1960) and compared with re-cycling three-time series (2) of instantaneous velocity planes at the input. These profiles were taken from a precursor simulation and a physical constraint were that of plane Couette flow at the step corner and the upwind boundary condition was used in previous numerical simulations (see Lowery and Reynolds (2007)). This gives an expression in which the Navier–Stokes equations are discretized in non-dimensional form:

\[ u, \hat{v}, \hat{w}, \hat{p} = \frac{u}{U_w}, \frac{v}{U_w}, \frac{w}{U_w}, \frac{p}{ho U_w^2} \]

1. Numerical approach

The computational domain has a length of \(L_x \), the expansion ratio of the two channels is defined by a single section \(L_x \), \(L_y \), \(L_z \), \(H \) in the normal direction and, \(U_0 \) in the quasisteady-state velocity. A total of \(256 \times 256 \times 128 \) grid points are used in the plane Couette flow as the large-scale motion, especially for a BFS configuration, where the flow structures of the step is the same as that stored by Reich (1978) in a fully-developed turbulent Couette flow.

2. Flow configuration and governing equations

Fig. 1 shows a schematic view of the Couette–backface–mixing-layer (CBL) system which is composed of a step of height \(H \) and an upstream channel heights, i.e. \(h \) and \(h_o \). By varying \(H /h_0 \), \(H /C_0 \), \(H /C_1 \), \(H /C_2 \), \(H /C_3 \), and \(H /C_4 \) is equal to that of the channel upstream, i.e. \(h /h_o \). This gives an expression in which the Navier–Stokes equations are discretized in non-dimensional form:

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turbulence. Patches of positive fluid velocity entering the incoming flow separates at the sharp step edge and reattaches further downstream leading to the formation of a primary recirculation region. The simulations were started from an arbitrary flow field and therefore irreversible to a statistically steady state. The time step used was \( \Delta t = 0.0033 \). Statistics were gathered for \( 300 \Delta t \), after the flow field had evolved into a statistically steady state.

In Fig. 5, we display a top view of iso-surfaces of turbulent plane Couette flow over a backward-facing step. The various plots of the velocity and vorticity fields presented in this section have been non-dimensionalized by \( \text{U}_w \) and \( \text{H} \), respectively.

4.1. Velocity

A snapshot of the flow field is shown in Fig. 1, where a three-dimensional slice of the simulation of the instantaneous velocities is presented. In order to compare our primary simulations with the data of Bech et al. (1995), the adopted grid spacing turned out to be too coarse. The simulations were started from an arbitrary flow field and therefore irreversible to a statistically steady state.

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structures grow with increased downstream propagation of the step and remain distributed over the entire domain. A roll-up is identified as a connected region of negative $\omega_x$ whereas in the secondary recirculation region the flow shows surprisingly strong spanwise motions.

The instantaneous velocity fluctuations do not exhibit the same streamwise coherence as the instantaneous vorticity fluctuations.

### 4.2 Vorticity

#### 4.2.1 Instantaneous Vorticity Structures

The instantaneous vorticity is decomposed into a mean vorticity component and vorticity fluctuations. The mean vorticity component is obtained by decomposing the velocity field into a mean and fluctuating part.

#### 4.2.2 Vortex Identification

These fluctuations do not exhibit the same streamwise coherence as the instantaneous vorticity fluctuations.

### 4.3 Vorticity

#### 4.3.1 Instantaneous Vorticity Contours

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### 4.4 Vorticity

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These fluctuations do not exhibit the same streamwise coherence as the instantaneous vorticity fluctuations.
5. Mean flow and turbulence statistics

The BFS simulation was performed at $Re_{x}=5200$ which led to $Re_{x}=140$ based on the step height and the upstream friction velocity $u_{f}$. Some comparisons will be made with the hot-wire data of Morinishi (2007) at $Re_{x}=300$ and expansion ratio 1.2.

5.1. Mean statistics

From pressure-driven BFS flows it is known that the upstream behaviour of the flow type is the most crucial due to the formation of a large primary separation bubble which extends about 7.5 times the step height at the step. A secondary separation bubble at $z_{h}/C_{0}$ is observed close to the primary separation bubble with clockwise motion observed in Fig. 10. An almost linear variation of $C_{f}$ is observed for the two cases. This implies that the streamwise-mean pressure gradient has become independent of $z$, in keeping with the measurements of Morinishi (2007), and the flow field can be considered as being nearly fully-developed in the downstream part of the computational domain. This is also consistent with the constancy of $C_{f}$ observed in Fig. 11.

The mean velocity profiles are presented in Fig. 12 and 13 for the mean and turbulent components at different upstream friction velocities $Re_{x}$ and expansion ratio $e$. The streamwise mean pressure gradient has become independent of $z$, in keeping with the measurements of Morinishi (2007), and the flow field can be considered as being nearly fully-developed in the downstream part of the computational domain. This is also consistent with the constancy of $C_{f}$ observed in Fig. 11.

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Although the final statistics in the lee of the step is relatively calm, the flow is by no means laminar.

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motion. Although the characteristic-S shape of the mean velocity profile \( U(x,y) \) has been retained at \( x \geq 25 \), the profile is yet to form being anti-symmetric. Although between the wall it is still wider at the upper half of the channel than at the lower half, no sign of being a fully developed Couette flow. However, irrespective of the length of the domain that can be used in the downstream part of the channel, an anti-symmetric profile corresponding to a fully developed Couette flow will not be reached. This is due to the principle of mass conservation. It follows that since the height of the domain after the step is twice that of the inlet section and the length of the domain that can be used in the downstream part of the channel is roughly half of the length of the domain after the step, the flow cannot adjust itself to an anti-symmetric profile since \( U(w) = 0 \) is a constant flow rate.

An examination of the mean wall-normal velocity profiles \( U_w(x,y) \) presented in Fig. 13 shows that the flow emanating from the step corner exhibits an upwards motion in the lower half of the channel and a downwards flow above \( y < 1 \). The distinct positive \( U_w(x,y) > 0 \) in the primary recirculation region is only obtained near the step corner in the lower half of the channel. This pattern is maintained up to \( y = 16 \), where after the turbulence levels of the spanwise and wall-normal components become dominant, the asymmetry in \( U_w(x,y) \) persists, where a substantially higher longitudinal turbulence intensity is observed near the moving wall that is almost twice that near the stationary wall. Apart from the secondary recirculation region, where \( U_w(x,y) = 0 \), the turbulence levels encompass the entire upper half of the channel, the asymmetry remains to be a most significant. The flow anisotropy was not accessible with the I-type hot-wire used in Morinishi’s (2007) measurements since only \( U_w(y) \) could be obtained.

The profiles shown in Fig. 14 show that the Reynolds shear stress is positive almost throughout the whole domain. The distinct positive \( \tau_{xy}(x,y) > 0 \) in the primary recirculation region is only obtained near the step corner in the lower half of the channel. Beyond the reattachment region, the mean wall-normal velocity is almost zero. Although the characteristic-S shape of the mean velocity profile \( U(x,y) \) has been retained at \( x \geq 25 \), the profile is yet to form being anti-symmetric. Although between the wall it is still wider at the upper half of the channel than at the lower half, no sign of being a fully developed Couette flow. However, irrespective of the length of the domain that can be used in the downstream part of the channel, an anti-symmetric profile corresponding to a fully developed Couette flow will not be reached. This is due to the principle of mass conservation. It follows that since the height of the domain after the step is twice that of the inlet section and the length of the domain that can be used in the downstream part of the channel is roughly half of the length of the domain after the step, the flow cannot adjust itself to an anti-symmetric profile since \( U(w) = 0 \) is a constant flow rate.

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that the contribution of the reattachment point at about 7 step heights downstream of the step. Further downstream, the profile fails to attain the symmetrical shape which characterizes a pure Couette flow; see e.g. Ref. 16 (1996). The profile of $u^+$ at $x = 17$, $22$ step heights downstream of the step, shows that the Reynolds shear stress is substantially higher near the moving wall than along the fixed wall.

2.2. Reynolds stress budget

In this section, the Reynolds stress budget is presented at a distance from the step and in a wall-normal region confined to the shear-layer region ($x < -1.5$). The transport equation for the Reynolds stress tensor is:

$$
\frac{D}{Dt} \rho \overline{u_i u_j} = -\overline{u_i \partial u_j / \partial x_k} \rho \overline{u_k} + \rho \overline{u_i} \overline{u_j} + \rho \overline{u_i} \overline{u_j} - \tau_{ij} - \rho \overline{u_i} \overline{u_j} + \rho \overline{u_i} \overline{u_j} - \rho \overline{u_i} \overline{u_j} + \rho \overline{u_i} \overline{u_j}
$$

where the production, dissipation, pressure-strain, pressure diffusion, molecular diffusion and turbulence diffusion are defined as:

- Production:
  $$
  \rho \overline{u_i} \overline{u_j} = -\rho \overline{u_i \partial u_j / \partial x_k} \overline{u_k}
  $$

- Dissipation:
  $$
  \rho \overline{u_i} \overline{u_j} = -\rho \overline{u_i} \overline{u_j}
  $$

- Pressure-strain:
  $$
  \rho \overline{u_i} \overline{u_j} = \frac{1}{2} \rho \overline{u_i} \overline{u_j}
  $$

- Pressure diffusion:
  $$
  \rho \overline{u_i} \overline{u_j} = \frac{1}{2} \rho \overline{u_i} \overline{u_j}
  $$

- Molecular diffusion:
  $$
  \rho \overline{u_i} \overline{u_j} = -\rho \overline{u_i} \overline{u_j}
  $$

- Turbulence diffusion:
  $$
  \rho \overline{u_i} \overline{u_j} = -\rho \overline{u_i} \overline{u_j}
  $$

The budget for $u^+$ in Fig. 16b is largely dominated by production and pressure-strain, where a large peak of negative production is observed close to the position of maximum $u^+$. Meanwhile, the profile of the molecular diffusion is mainly positive in the near-wall turbulent flow.

3.2. Recovery region

In this section, the Reynolds stress budget is presented at a distance from the step and in a wall-normal region confined to the shear-layer region ($x > 10$). The transport equation for the Reynolds stress tensor is:

$$
\frac{D}{Dt} \rho \overline{u_i u_j} = -\overline{u_i \partial u_j / \partial x_k} \rho \overline{u_k} + \rho \overline{u_i} \overline{u_j} + \rho \overline{u_i} \overline{u_j} - \tau_{ij} - \rho \overline{u_i} \overline{u_j} + \rho \overline{u_i} \overline{u_j} - \rho \overline{u_i} \overline{u_j} + \rho \overline{u_i} \overline{u_j}
$$

where the production, dissipation, pressure-strain, pressure diffusion, molecular diffusion and turbulence diffusion are defined as:

- Production:
  $$
  \rho \overline{u_i} \overline{u_j} = -\rho \overline{u_i \partial u_j / \partial x_k} \overline{u_k}
  $$

- Dissipation:
  $$
  \rho \overline{u_i} \overline{u_j} = -\rho \overline{u_i} \overline{u_j}
  $$

- Pressure-strain:
  $$
  \rho \overline{u_i} \overline{u_j} = \frac{1}{2} \rho \overline{u_i} \overline{u_j}
  $$

- Pressure diffusion:
  $$
  \rho \overline{u_i} \overline{u_j} = \frac{1}{2} \rho \overline{u_i} \overline{u_j}
  $$

- Molecular diffusion:
  $$
  \rho \overline{u_i} \overline{u_j} = -\rho \overline{u_i} \overline{u_j}
  $$

- Turbulence diffusion:
  $$
  \rho \overline{u_i} \overline{u_j} = -\rho \overline{u_i} \overline{u_j}
  $$

The budget for $u^+$ in Fig. 16c is largely dominated by production and pressure-strain, where a large peak of positive production is observed close to the position of maximum $u^+$. Meanwhile, the profile of the molecular diffusion is mainly positive in the near-wall turbulent flow.
Field data were considered by these authors in which the mean-removed data was varied by changing the wall velocity and the streamwise pressure gradient of particular relevance to our study is case CII of Kuroda et al. (1995).

Owing to the asymmetry of the Couette–Poiseuille flow, the wall shear stresses on the lower and upper-wall are generally different. For the current simulation, the corresponding local wall shear-stresses are predominantly aligned in the streamwise direction, even though the fluctuating vorticity field, associated with the wall-normal turbulence, is close to isotropy.

The rather different behaviour in the vicinity of the lower-wall is due to the nearly vanishing friction velocity $u_{10}$. In the present case, however, the results show that the vorticity production only plays a marginal role in the lower-wall.

The mean velocity profile and turbulent intensities presented in Fig. 17 compare surprisingly well with the DNS data of Kuroda et al. (1995), especially near the upper-wall. Substantial differences are observed near the lower-wall. The streamwise turbulence intensity, for instance, exhibits a modest peak in the present case, whereas a maximum intensity occurs near the wall. The large-scale turbulence, is close to isotropy.

The wall-normal vorticity fluctuations inevitably vanish at both walls due to the no-slip conditions. The local near-wall peak of the wall-normal $\alpha = \partial \omega_z / \partial x_2$ is associated with streamwise-oriented vortices, which are amplified and attaining higher values than $\alpha = \partial \omega_z / \partial x_1$. This indicates that the turbulent vortices are predominantly aligned in the streamwise direction, even though the fluctuating vorticity field, associated with the wall-normal turbulence, is close to isotropy.

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The symbols denote DNS data from Kuroda et al. (1995).

Fig. 18. Wall-normal distribution of the root-mean-square vorticity fluctuations normalized by $k$. (a) $x/h = 16$ and (b) $x/h = 18$.

Fig. 19. Wall-normal distribution of the root-mean-square vorticity fluctuations normalized by $k$. (a) $x/h = 16$ and (b) $x/h = 18$.
A direct numerical simulation of turbulent Couette flow over a backward-facing step has been performed at a relatively low Reynolds number. The mean reattachment length of the Dean vortices was found to be 1.7. In the recirculation zone a large negative skin friction coefficient was observed beneath the core of the primary separation bubble.

The nonlinear finite element structure formed along the step perpendicular to the step height downstream of the step, and the reattachment point was found to be downstream of the step. High and anisotropic turbulence levels were produced in the shear-layer downstream of the step. The shear-layer which formed when the upstream Couette flow turned towards the reattachment point and the shear-layer is sustained several step heights downstream of the step. High and anisotropic turbulence was found to be present in the shear-layer downstream of the step. The turbulence in the approaching Couette flow is caused by the primary separation bubble, which gives rise to a large negative skin friction coefficient beneath the core of the primary separation bubble.

Even though the wall-normal mean velocity $V$ is almost zero at downstream of the step, the shear-layer's velocity did not retain the characteristic $T$-shape typical of a pure Couette flow. This phenomenon is attributed to the principle of global mass conservation, which can be fulfilled only if an adverse mean pressure gradient is present underneath the core of the primary separation bubble. In the recirculating asymmetric mean flow, the shear-layer's velocity was roughly 10 times higher along the moving wall than at the steady recirculation. The resulting flow field was thus far less energetic along the stagnant wall and was reflected in the low levels of velocity and vorticity fluctuations.

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