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On dynamics and secondary currents in meandering confined turbulent shallow jet

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ABSTRACT

We study a shallow turbulent confined jet issuing from a wide duct of the aspect ratio around 0.1 ($\approx H : B =$ height: width) into a large rectangular reservoir ($200H \times 267H$ long in spanwise and streamwise direction) bounded by two parallel plates with the distance H using a wall-resolved Large-eddy simulation. According to statistical and dynamical features of the flow the domain is divided into the near ($< 40H \approx 4B$), middle and far ($> 130H \approx 13B$) field along the streamwise coordinate. The well-known meandering (or flapping) motion of the jet core starts at $\approx 5B$ generating large-scale planar vortices similar to the von Kármán vortex street. We find that the typical smaller scale coherent structure appears to be a vortex tube (with the cross-section diameter $< H/2$) oriented along the flow with a high value of the streamwise vorticity. In the near field a pair of these counter-rotating vortices are produced in each shear layer with the help of the strong entrainment of the ambient fluid into the turbulent jet core and look similar to the braid vortices connecting Kelvin–Helmholtz rolls found in the classical mixing layer. The maximum time-averaged streamwise vorticity is reached at the axial distance of $12.5H$ with monotonic decrease further downstream. In the middle and far fields these streamwise rolls are mostly observed in one shear layer (on one side of the jet) while being suppressed on the other side by the flapping motion. Typically they are organized in a "zig-zag" chain across the "shallow" direction. A simple model is proposed to explain this feature.

1. Introduction

Wall-bounded turbulence is populated with coherent vortices which significantly contribute to statistical flow characteristics. A specific class of coherent structures appears in a flow issuing from a slit into a reservoir bounded by two narrow parallel plates (slot jet). Such types of flows are usually referred to as "quasi-two-dimensional" since spanwise and streamwise directions are much larger than the distance between rigid surfaces (Jirka, 2001; Jirka and Uijttewaal, 2004; Heijst and Clercx, 2009; Uijttewaal, 2014).

At certain shallow flow parameters a well-known meandering pattern (Giger et al., 1991; Dracos et al., 1992; Chen and Jirka, 1999; Landel et al., 2012a; Shestakov et al., 2015) appears as a consequence of a sinous linear instability (Chen and Jirka, 1998). The meandering motion induces planar counter-rotating large-scale vortices squeezed between bounding plates. According to Landel et al. (2012a), they propagate downstream approximately four times slower than the jet core and settle in both shear layers in a checkerboard order similar to the classic von Kármán vortex street. This flow regime decays and ceases to exist at certain location downstream from the nozzle exit as a consequence of the wall friction significantly destroying the momentum of the jet. Recently, we noticed that these dominant vortical structures may detach from pressure eddies contradicting to typical pressure-based vortex eduction criteria (Hanjalić and Mullyadzhanov, 2015).

The near field region of a shallow wall-bounded jet is governed by secondary currents due to interactions of mixing layers and walls (Holdeman and Foss, 1975; Rockwell, 1977). Since there is a considerable decrease of the normal-to-the-wall velocity fluctuations it has been argued that far field provides only two-dimensional flow motions (Giger et al., 1991; Dracos et al., 1992). The usual practice to model "shallow" direction is to add the "Rayleigh friction" term (linear friction) to the 2D Navier–Stokes equation (Heijst and Clercx, 2009). However, possible three-dimensional flow currents distort a Poiseuille-like velocity

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profile questioning this modeling approach (Akkermans et al., 2008a).

The influence of a "shallow" direction on three-dimensional character of the flow is usually addressed in the context of electromagnetically forced flows in a bounded container (Lardeau et al., 2008; Shats et al., 2010; Kelley and Ouellette, 2011; Xia et al., 2011; Akkermans et al., 2012) and an evolution of dipole vortices (Lin et al., 2003; Sous et al., 2004, 2005; Akkermans et al., 2008a; Akkermans et al., 2008b; Duran-Matute et al., 2010). Recent experiments by Xia et al. (2011) describe the interaction of a large-scale horizontal vortex and a small-scale normal-to-the-wall turbulent motion. The latter is suppressed, which leads to an inverse energy cascade. In case of the evolution of an impulsively started quasi-two-dimensional dipole, a relatively strong circulation across the thin fluid layer is observed on the front side of the vortex (Lin et al., 2003; Sous et al., 2004, 2005; Akkermans et al., 2008a,b). To our knowledge no systematic three-dimensional analysis has been performed for spatially evolving stationary (as opposed to decaying dipole vortices) shallow jets, and here we report on new physical phenomena observed in these advection-dominated flows.

In this paper we perform a three-dimensional Large-eddy simulation (LES) of a shallow turbulent wall-bounded jet. The focus is on secondary currents between narrow plates and their evolution downstream. The inflow is provided with a fully developed turbulent duct flow of the aspect ratio about $B : H \approx 10$, where B is the width of the incoming channel and H is a small distance between narrow parallel plates. We believe that the shallow jet issuing from a "wide" source as is considered in the present study is natural for environmental applications (see Fig. 1 in Giger et al. (1991), Fig. 1 and discussion in Rowland et al. (2009)). The flow in the present geometry also provides a clear classification of the near-, middle- and far-field of the jet as is discussed below. The Reynolds number is 20×10^3 based on $H/2$ and the bulk velocity W_0 of the duct flow. We note that for the present geometry parameters the flow loses stability around $5B$ downstream and starts flapping in the spanwise direction. This meandering motion generates large-scale vortices that grow in the shear layers. These structures are easy to detect and are the central topic of many studies. We show that the typical small-scale coherent structure is a streamwise-oriented vortex roll both in the near field where mixing layers do not influence each other and in the middle and far field where these structures meander together with the main flow.

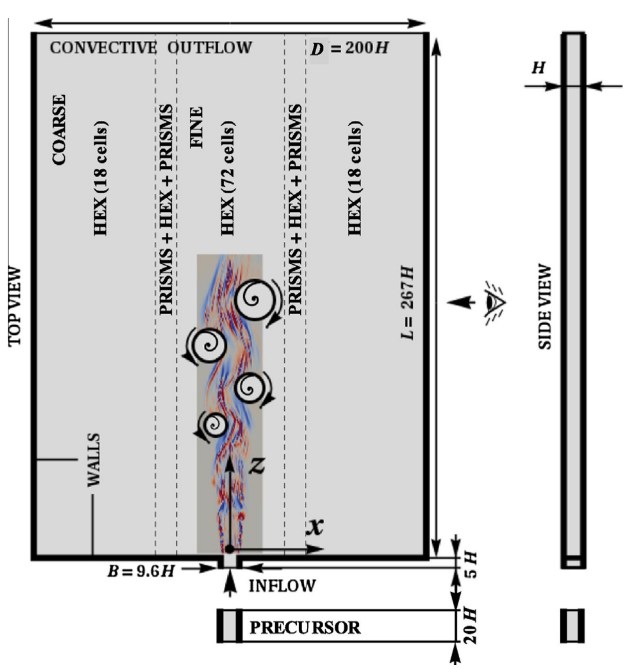

Fig. 1. The domain considered for numerical simulations. The coordinate system is placed at the nozzle exit. Two narrow plates correspond to $y = 0$ and $y = H$.

2. Governing equations and numerical details

We solve the spatially filtered momentum and continuity (LES) equations to determine the dynamics of the fluid:

$$\frac{\partial U_i}{\partial t} + \frac{\partial U_i U_j}{\partial x_j} = -\frac{1}{\rho}\frac{\partial P}{\partial x_i} + \frac{\partial}{\partial x_j}\left[(\nu + \nu_t) 2 S_{ij}\right],$$

$$\frac{\partial U_i}{\partial x_i} = 0,$$

where ρ, U_i, P and ν are the constant density of the fluid, components of the filtered velocity vector $\mathbf{U} = (U, V, W)$, pressure field and kinematic viscosity, repsectively. The sub-grid scales are closed with the eddy-viscosity dynamic Smagorinsky model with $\nu_t = (C_s \Delta)^2 \sqrt{2 S_{ij} S_{ij}}$, where $S_{ij} = (\partial U_i/\partial x_j + \partial U_j/\partial x_i)/2$ is the rate-of-strain tensor. The Smagorinsky constant C_s is determined with the standard dynamic procedure (Germano et al., 1991), while the characteristic length Δ is the local grid size.

The LES is performed using the finite-volume unstructured computational code T-FlowS (Ničeno and Hanjalić, 2005), with the cell-centered collocation grid arrangement. The diffusion and convection terms in the momentum equations are discretized by the second-order central-differencing scheme, while a fully implicit three-level time scheme is used for time-marching. The velocity and pressure are coupled by the iterative pressure correction algorithm SIMPLE.

We consider a wall-bounded shallow jet flow in the rectangular domain of the size $D \times H \times L$ ($x \times y \times z$), where D ($= 200H$) and L ($= 267H$) are much larger than H (see Fig. 1). The size of the domain is taken in accord with the experimental facility at the Institute of Thermophysics to mimic the ongoing measurements (Shestakov et al., 2015). The inflow channel is a short duct of the length $5H$ (and the cross-section $B \times H$ with $B = 9.6H$), which is meshed with $N_x \times N_y \times N_z = 142 \times 72 \times 25$ hexahedral cells in spanwise (x), wall-normal (y), and streamwise (z) directions, respectively. The inflow velocity profiles are taken from a precursor simulation of a fully developed turbulent duct flow with dimensions $B \times H \times 20H$ and periodic streamwise boundary conditions (see the bottom of Fig. 1), which is covered with $N_x \times N_y \times N_z = 142 \times 72 \times 100$ hexahedral cells. Two bands between $27 < |x|/H < 33$ in the receiving volume contain two layers of triangular prisms and one layer of hexahedral cells in between to reduce the streamwise and wall-normal number of cells in the rest of the domain. It allows us to use 18 cells in y instead of 72 and 350 cells in z direction instead of 720. The resulting grid consists of 23.74 mln cells. The convective outflow condition is used at the end of the domain, while the no-slip condition is applied for all walls. The grid points are clustered towards the walls and within the mixing layers with an increment less than 5 percent.

To check the convergence of the obtained LES solution we compare time-averaged (denoted by overbar) centerline axial velocity profiles from the present and previous coarser computation with about 15 mln cells (Mullyadzhanov et al., 2015). Note that according to Fig. 2 the comparison shows that two profiles for $z/H < 130$ are in very good agreement while further downstream the discrepancy becomes more obvious. In the relevant literature the jet is usually divided into near, middle and far fields. The region after $z/H \approx 130$ ($z \approx 13B$) we call the far field as the centerline time-averaged axial velocity reaches $z^{-1/2}$ dependence in agreement