Polydisperse turbidity currents propagating over complex topography: Comparison of experimental and depth-resolved simulation results

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A computational investigation is presented of mono-, bi-, and polydisperse lock-exchange turbidity currents interacting with complex bottom topography. Simulation results obtained with the software TURBINS are compared with laboratory experiments of other authors. Several features of the flow, such as deposit profiles, front location, suspended mass, and runout length, are discussed. For a monodisperse lock-exchange current propagating over a flat surface, we investigate the influence of the boundary conditions at the streamwise and top boundaries, and we generally find good agreement with corresponding laboratory experiments. However, we note some differences with a second set of experimental data for polydisperse turbidity currents over flat surfaces. A comparison with experimental data for bidisperse currents with varying mass fractions of coarse and fine particles yields good agreement for all cases except those where the current consists almost exclusively of fine particles. For polydisperse currents over a two-dimensional bottom topography, significant discrepancies are observed. Possible reasons are discussed, including erosion and bed load transport. Finally, we investigate the influence of a three-dimensional Gaussian bump on the deposit pattern of a bidisperse current.

1. Introduction

Turbidity currents represent an important mechanism by which sediment is transported from the continental shelves into the deep oceans (Meiburg and Kneller, 2010). They play a crucial role within the global sediment cycle and, over geological time scales, in the formation of hydrocarbon reservoirs (Syvitski et al., 1996). Turbidity currents can reach front velocities \( O(10) \) m/s, travel up to \( O(1,000) \) km, and transport \( O(100) \) km\(^3\) of sediment.

The formation of topographical features such as channels, gullies, sediment waves, and levees can occur as a result of the depositional and/or erosional nature of turbidity currents interacting with the seafloor (Wynn et al., 2000a; Migeon et al., 2001; Nakajima and Satoh, 2001; Wynn and Stow, 2002). Conversely, the seafloor topography can influence the dynamics of turbidity currents, and hence their deposits. Examples of such effects include the interaction of turbidity currents with large barriers (Pickering and Hiscott, 1985; Kneller and McCaffrey, 1999), differential deposition due to slope gradients (Garcia and Parker, 1989; Kubo, 2004), and flows in meandering channels (Imran et al., 1999; Straub, 2007).

Due to the unpredictable and often catastrophic nature of turbidity currents, there has been an intensive effort to study their behavior by means of laboratory experiments (e.g., Luthi, 1981a; Bonnecaze et al., 1993; Gladstone et al., 1998; de Rooij and Dalziel, 2001), theoretical approaches including box models and shallow water theory (Luthi, 1981b; Rotman and Simpson, 1983; Bonnecaze et al., 1993; Baines, 1995; Dade and Huppert, 1995; Hallworth et al., 1998), and more recently, via numerical simulations (Necker et al., 2002, 2005; Kassem and Imran, 2004; Blanchette et al., 2005; Huang et al., 2008; Nasr-Azadani and Meiburg, 2011).

These investigations frequently address the so-called lock-exchange currents; cf. Fig. 1. In the current study, we employ the simulation tool TURBINS to address such flows, which was recently developed in our group (Nasr-Azadani and Meiburg, 2011). TURBINS is a highly parallel code that employs the Navier–Stokes equations in the Boussinesq approximation to describe the motion of the suspension, along with a transport equation for the particle concentration field. We employ TURBINS to reproduce and compare with experiments and numerical simulations by Gladstone et al. (1998), de Rooij and Dalziel (2001), Necker et al. (2002), and Kubo (2004). All of these studies focus on the depositional behavior of lock-exchange, mono-, and/or polydisperse currents flowing over...
flat or complex geometries. Specifically, we will focus on such quantities as front location, runout length, and deposit profiles of fine and coarse particles. Of particular interest are the coupling mechanisms between particles of different sizes, and their effects on the transient and final deposit profiles. In addition, we investigate the deposition of a bidisperse suspension interacting with a three-dimensional Gaussian bump, cf. Fig. 1. This will shed light on the influence of complex bottom topography with regard to generating nonuniformities in the final deposit profiles. The emphasis of the current investigation is not so much on the highly complex topographies that might be encountered on a real world sea floor. Rather, we focus on geometrically simple shapes such as the ones that have been employed in various laboratory experiments. The hope is that the insight we gain from studying the interaction of turbidity currents with such relatively simple geometries will allow us to gain an understanding of the fundamental mechanisms that may dominate the interaction between currents and realistic sea floor topographies. Hence we aim to address such important issues as the relationship between nonuniformities in the deposit profiles and the shape of the obstacle, the effect of the geometry on the transport and deposition of fine and coarse particles, and the runout length of these currents.

Section 2 formulates the governing equations. A brief discussion of the numerical methodology is presented in Section 3. Section 4 presents the numerical results obtained via TURBINS for a monodisperse lock-exchange flow, and validates them against experiments conducted by de Rooij and Dalziel (2001). Subsequently, we proceed towards a comparison with experimental results for bidisperse turbidity currents (Gladstone et al., 1998). We analyze bottom shear stress data to evaluate the potential for bed load transport and erosion on the final deposit profiles. Finally, we investigate the influence of the bottom topography on nonuniformities in the deposits of polydisperse suspensions.

2. Governing equations

We employ the incompressible Navier–Stokes equations for the motion of the suspension, along with a convective–diffusive transport equation for the Eulerian description of the particle concentration field. A detailed discussion of the governing equations is provided by Necker et al. (2002) and Nasr-Azadani and...
Meiburg (2011), so it suffices to present a brief summary here.

We focus on modeling incompressible flows of dilute suspensions, with typical particle volume fractions of $O(1\%)$ or less. The suspension flow is governed by the incompressible Navier–Stokes equations in the Boussinesq approximation:

$$\nabla \cdot \mathbf{u} = 0,$$

(1)

$$\frac{\partial \mathbf{u}}{\partial t} + \mathbf{u} \cdot \nabla \mathbf{u} = -\nabla p + \frac{1}{\mathrm{Re}} \nabla^2 \mathbf{u} + \mathbf{C}_1 \mathbf{e}^s,$$

(2)

where the nondimensional quantities $\mathbf{u}$, $p$, $\mathrm{Re}$, and $\mathbf{C}_1$ represent the fluid velocity, pressure, Reynolds number, and total particle concentration (volume fraction), respectively. $\mathbf{e}^s$ indicates the unit vector in the direction of gravity, and $\mathbf{C}_1 = (0, -1, 0)$. By assuming a dilute suspension of small particles, we can neglect particle inertia and any particle–particle interaction. Consequently, the particles are assumed to follow the fluid, with a constant settling speed $u_s$ superimposed that acts in the direction of gravity. We allow the suspension to contain $N$ different particle sizes, all of the same density $\rho_p$. For the $i$th particle size, we define a continuum concentration field $\mathbf{C}_i(x, t)$, which is governed by

$$\frac{\partial \mathbf{C}_i}{\partial t} + \mathbf{u} \cdot \nabla \mathbf{C}_i = \frac{1}{\mathrm{Re} \, \mathbf{S}_c} \nabla^2 \mathbf{C}_i, \quad i = 1, \ldots, N.$$

(3)

Here, $u_s^i$ and $\mathbf{S}_c$ denote the $i$th particle settling speed and Schmidt number

$$\mathbf{S}_c = \frac{\nabla}{K_1}, \quad i = 1, \ldots, N,$$

(4)

respectively. In Eq. (4), $\nabla$ and $K_1$ represent the kinematic viscosity of the fluid and the diffusion coefficient associated with the ith concentration field. We remark that in Eq. (3), diffusion is included to account for the smearing of the sharp interfaces of particles (Ham and Homsy, 1988; Davis and Hassen, 1988) due to hydrodynamic diffusion of particles. Hertel et al. (2000b) demonstrate that this diffusion does not significantly affect the dynamics of the flow, as long as the Schmidt number is equal to or greater than one. In the present investigation, we solve all $\mathbf{S}_c$’s to unity. Unless stated otherwise, for the given particle size, the settling velocity is taken to be the Stokes settling velocity. For small, spherical particles with low particle Reynolds number, the Stokes settling velocity agrees closely with that of the empirical correlations (Gibbs et al., 1971; Dietrich, 1982). On the other hand, for irregularly shaped particles, considerable deviation from the Stokes value may occur in the settling velocity, cf. Gladstone et al. (1998). All concentration fields $\mathbf{C}_i$ are scaled with the total initial volume fraction of the particles in the lock $\mathbf{C}_i$; i.e.,

$$\mathbf{C}_i = \frac{\mathbf{C}_i}{\mathbf{C}}, \quad i = 1, \ldots, N.$$

(5)

The total concentration $\mathbf{C}_i$ introduced in the momentum equation (see Eq. (2)) is obtained as

$$\mathbf{C}_i = \sum_{i=1}^{N} \mathbf{C}_i.$$

(6)

Here, $\mathbf{C}_i$ varies between 0 in the clear fluid, and 1 in the lock (see Fig. 1). In the remainder, we choose half the lock height as the characteristic length scale, i.e. $L_y = H/2$. The buoyancy velocity

$$u_b = \sqrt{\frac{\rho_g - \rho_0}{2 \rho_0}} \nabla \mathbf{H},$$

(7)

serves as the characteristic velocity to nondimensionalize the flow quantities. Consequently, the Reynolds number (see Eqs. (2) and (3)) is defined as

$$\mathrm{Re} = \frac{u_b H/2}{\nu}.$$  

(8)

Note that the * symbol refers to dimensional quantities.

### 3. Numerical method

A detailed description of the numerical method is presented in Nasr-Azadani and Meiburg (2011). We solve the momentum equations using a projection method (Chorin, 1968) in conjunction with the fractional step method (Kim and Moin, 1985), on a MAC-staggered grid. In this approach, the velocity components are stored at the cell faces, while scalar quantities including pressure and concentration(s) are stored at the cell centers. Viscous and diffusion terms in the momentum and transport equations are discretized via a fully implicit central differencing method, when the convective terms are discretized via an explicit third-order essentially nonoscillatory (ENO) scheme (Harten et al., 1987). The time integration for the transport and momentum equations is performed via a third-order total variation diminishing Runge–Kutta method (TVD-RK3; cf. Harten, 1997). This method ensures a solution without any spurious oscillations, as long as no oscillations are generated by the Euler method during each substep integration. We briefly present one Runge–Kutta substep (not the entire step) of the time integration procedure in the following. First, the concentration fields of all particle sizes are updated from time level $t_n$ to $t_{n+1}$ as

$$\frac{\hat{c}_i^{n+1} - c_i^n}{\Delta t} + (u^n + u'e^s) \cdot \nabla c_i^n = \frac{1}{\mathbf{S}_c \, \mathrm{Re}} \nabla^2 c_i^{n+1} + f_i^{n+1},$$

(9)

where $f_i$ is the forcing term employed to apply the boundary conditions on the solid boundaries (Nasr-Azadani and Meiburg, 2011). Next, the velocity field is updated to an intermediate velocity field $\mathbf{u}^*$

$$\frac{\mathbf{u}^* - \mathbf{u}^n}{\Delta t} + \nabla \mathbf{u}^n = -\nabla p^n + \frac{1}{\mathrm{Re}} \nabla^2 \mathbf{u}^* + \mathbf{C}_1^{n+1} + f^*,$$

(10)

with $f^*$ accounting for the forcing term for boundary condition imposition on the solid boundaries. The divergence-free velocity field $\mathbf{u}^{n+1}$ at the new time level $t_{n+1}$ is obtained via

$$\frac{\mathbf{u}^{n+1} - \mathbf{u}^*}{\Delta t} = -\nabla \psi^{n+1},$$

(11)

where the scalar variable $\psi$ is computed from the solution of a Poisson equation of the form

$$\nabla^2 \psi^{n+1} = \frac{1}{\Delta x} \nabla \cdot \mathbf{u}^*.$$  

(12)

Again, we remark that Eqs. (9)–(12) represent only one out of the three Runge–Kutta substeps associated with the TVD-RK3 methodology. We refer the reader to Nasr-Azadani and Meiburg (2011) for a detailed description of the solution procedure.

We have adopted an immersed boundary method with direct forcing (Mittal et al., 2008) to enforce the boundary conditions accurately on the solid surface. This method benefits from ease of implementation, no additional constraints on the time step, and accurate calculation of shear stress profiles on the solid surface. The last is of great importance when erosion of particles from the sediment bed into the flow plays a significant role.

Fig. 2 shows the sketch of the Cartesian grid, along with different nodes tagged in accordance with the location of the solid surface $I$. We employ a bilinear (trilinear) in two (three) dimensions interpolation scheme to obtain the value of any fluid quantity $q$ at the mirrored node $O$, based on the surrounding fluid
The initial volume fraction of the particles is based on a lock-exchange configuration in which the suspension is dilute. The Stokes and focusing on particle-driven flows, in order to demonstrate the ability of TURBINS to accurately reproduce experimental and computational results provided by de Rooij and Dalziel (2001)

**4. Results**

4.1. Validation

Nasr-Azadani and Meiburg (2011) presented a detailed validation study of the numerical methodology implemented in TURBINS. Specifically, they were able to demonstrate the expected performance in terms of accuracy and convergence for constant density and gravity-driven flows over complex geometries. In this section, we will extend their investigation by focusing on particle-driven flows, in order to demonstrate the ability of TURBINS to accurately reproduce experimental and computational results provided by de Rooij and Dalziel (2001) and Necker et al. (2002). Specifically, we will focus on such quantities as front velocity, runout length, and deposit profiles.

The experiments carried out by de Rooij and Dalziel (2001) are based on a lock-exchange configuration in which the suspension is initially in a reservoir of \( L_x \times L_y \times H = 10 \times 26 \times 26.5 \) cm filled with tap water. The suspension includes 50 g of silicon carbide particles with diameter \( d_p = 37 \) \( \mu \text{m} \) and density \( \hat{\rho}_p = 3.217 \text{ g/cm}^3 \). The initial volume fraction of the particles is \( C_i = 0.0023 \), which allows us to assume that the suspension is dilute. The Stokes and Dietrich (1982) relationships result in identical particle settling velocities of \( \hat{u}_s = 0.16 \text{ cm/s} \). The reference length \( L_r = 13 \) cm yields a buoyancy velocity (see Eq. (7)) of \( \hat{u}_b = 8 \text{ cm/s} \). The nondimensional settling velocity and the Reynolds number (see Eq. (8)) hence take the values \( \hat{u}_s = 0.02 \) and \( Re = 10,400 \), respectively.

Fig. 3 shows the computational setup employed for our numerical simulation. The domain has a size of \( L_x \times L_y = 20 \times 2 \). HärTEL et al. (2000b) observe that for sufficiently large values of Reynolds number, such flow quantities as the front velocity depend only very weakly on \( Re \). Hence, we employ a value of \( Re = 5000 \) in our simulations, in order to limit the computational effort. We employ a uniform grid in both \( x \) - and \( y \)-directions, with grid spacings of \( \Delta x = 0.02 \) and \( \Delta y = 0.01 \), respectively. The lock, in which the fluid initially is at rest, has a dimensionless height of \( H = 2 \) and length of \( L_y = 0.75 \), respectively.

The appropriate set of boundary conditions for the solution of the concentration field includes no-flux boundary conditions at the left, right, and top walls, respectively:

\[
\frac{\partial \varepsilon_c}{\partial x} = 0 \quad x = 0, L_x, \quad (15)
\]

\[
\frac{1}{Sc \, Re} \frac{\partial \varepsilon_c}{\partial y} = 0 \quad y = L_y. \quad (16)
\]

At the bottom wall, the particles should leave the computational domain freely with a settling velocity \( \hat{u}_s \). This means that the effect of the evolving thin deposit layer on the subsequent flow is neglected. To this end, Necker et al. (2002) employed a convective outflow boundary condition for the concentration field. For the current MAC-staggered grid, we accomplish this in the following way. At the first concentration node above the bottom boundary, which is located half a grid cell above the bottom wall, we solve the usual transport equation (see Eq. (3)) for the concentration field. In doing so, we assume \( \partial \varepsilon_c / \partial y = 0 \) at \( y = 0 \). We remark that this approach is applicable only to situations where resuspension of particles from the bottom bed into the flow is negligible, which is the case for the experiments of de Rooij and Dalziel (2001). This was analyzed in detail by Necker et al. (2002), based on the approach of Shields (1936), Mantz (1973) and Yalin and Karahan (1979). Necker et al. (2002) demonstrated that, based on the maximum bottom shear stress observed in their simulations, resuspension of particles was unlikely to have occurred in the experiments of de Rooij and Dalziel (2001).

For the velocity field, we employ no-slip conditions at the bottom and right walls. Since the experiments were conducted in

![Fig. 2. Sketch of the solid surface \( \Gamma \) and the tagged nodes.](image)

**Table 1** Velocity boundary conditions for three different particle-driven lock-exchange simulations discussed in the text (see Fig. 3).

<table>
<thead>
<tr>
<th>Case</th>
<th>Left wall boundary condition</th>
<th>Top wall boundary condition</th>
</tr>
</thead>
<tbody>
<tr>
<td>a</td>
<td>Free-slip</td>
<td>No-slip</td>
</tr>
<tr>
<td>b</td>
<td>No-slip</td>
<td>Free-slip</td>
</tr>
<tr>
<td>c</td>
<td>No-slip</td>
<td>No-slip</td>
</tr>
</tbody>
</table>
a tank with a free surface, it would seem reasonable to apply a free-slip boundary condition at the top boundary of the computational domain. However, as discussed by Britter and Simpson (1978), a no-slip boundary condition may be more appropriate, as impurities in the water may cause the free surface to behave similarly to a solid wall. To investigate the effect of the conditions applied at both the top and left boundaries, we present results from three sets of simulations as described in Table 1. The deposited particle layer thickness can be obtained by integrating the particle flux through the bottom boundary in time:

$$D(x,t) = \frac{C_s}{\sigma} \int_0^t u_s c_w(x,t) \, dt.$$  \hspace{1cm} (17)

Here, $c_w$ denotes the concentration at the lower boundary. The constant $\sigma$ is introduced as the packing fraction of the settled sediment. It typically takes an approximate value of $\sigma = 0.63$ (Torquato et al., 2000). However, since the experiments referred to above report deposit profiles only in terms of particle mass, not layer thickness, we are not concerned with the porosity of the sediment layers in the present study, so that we simply set $\sigma = 1$.

Fig. 4 shows the final deposit profile along the bottom boundary for the experiment (de Rooij and Dalziel, 2001), along with corresponding two-dimensional simulation results for the three cases presented in Table 1. The overall agreement is good, with differences being primarily evident in the lock region. As discussed by Necker et al. (2002), possible reasons for the discrepancy can be found in the initial turbulence added to the suspension due to stirring (Gladstone et al., 1998), the initial suspension not being perfectly homogeneous, and perturbations introduced by the mechanical removal of the lock-gate. Finally, we cannot exclude the possibility that, in spite of the stirring applied, some particles may nevertheless have already settled out within the lock by the time the gate is opened, which may explain the relatively large amount of deposit within the lock region. The relatively strong oscillations observed in our results are not unexpected, as two-dimensional simulations at fairly low
Reynolds numbers show more coherent vortical structures than comparable three-dimensional simulations (Necker et al., 2002).

Fig. 5 displays the front location and the total suspended mass in the domain, as functions of time. Case (a) with free-slip boundary condition at the left wall shows good agreement with the results of Necker et al. (2002), who also employed a free-slip boundary condition at the left wall, due to the constraints on their numerical method.

A comparison of the results in Figs. 4 and 5 suggests that the condition at the top boundary does not have a significant influence on the front location and the final deposit profile. In contrast, the difference between implementing a no-slip or free-slip boundary condition at the left wall is more significant, which is a somewhat surprising result. We observe that in case (a), the current travels faster and maintains more particles over longer times. Also, the oscillatory behavior of the deposit profile near the lock-gate is less pronounced for this case, Fig. 4. In order to understand the reason behind the significant influence of the left wall boundary condition, we show instantaneous contours of the concentration field at different times in Fig. 6.

For the no-slip condition at the left wall, a packet of particles (indicated by the symbol ‘‘’’ in Fig. 6) detaches early on from the bulk of the current, due to the drag exerted by the left wall. These particles settle out fairly quickly, and hence are no longer available to drive the current but contribute to the increased deposited mass within the lock region for the no-slip case. We furthermore note that, regardless of the boundary condition imposed at the top wall, the vortical structure of the flow field is similar at corresponding times (see Fig. 6).

4.2. Bidisperse particle-driven current

Gladstone et al. (1998) investigated the effects of having a bidisperse mixture of particle sizes on the sedimentation behavior and runout length of turbidity currents. Toward this end, they varied the initial mass fractions of fine and coarse particles, while...
keeping the total volume fraction of particles in the lock at a constant value of \( C_v = 0.00349 \). Here, we duplicate their experiments by conducting two-dimensional simulations.

The lock used in the experiment has dimensions of \( L_v \times H \times W = 20 \times 40 \times 20 \) cm. The fine and coarse particles consist of silicon carbide with density \( \rho_p = 3.217 \) g/cm\(^3\), and they have diameters \( d_1 = 25 \) µm and \( d_2 = 69 \) µm, respectively. Gladstone et al. (1998) suggest that the particle settling velocity is approximately one-third smaller than the Stokes value, due to the additional drag originating from the angular particle shapes; cf. Hallermeier (1981) and Sparks et al. (1991). In this way, we obtain settling velocities of fine and coarse particles as \( \bar{u}_s^1 = 0.05 \) cm/s and \( \bar{u}_s^2 = 0.387 \) cm/s, respectively. Clearly, there is some uncertainty and possibly nonuniformity associated with the settling velocity, due to variations in the shapes of individual particles, and the mixture not being perfectly sorted.

With the reference length \( L_v = H / 2 = 20 \) cm and a particle volume fraction of \( C_v = 0.00349 \), the buoyancy velocity and corresponding Reynolds number are computed as \( \bar{u}_b = 12.33 \) cm/s and \( Re = 24,700 \), respectively; cf. Eqs. (7) and (8).

The computational setup is as follows (see Fig. 3). The domain has a size of \( L_z \times L_x = 32 \times 2 \), with uniform grid spacing (\( \Delta x = 0.032 \) and \( \Delta y = 0.01 \)) in both directions. Again, we employ a reduced value of \( Re = 4000 \) to limit the computational cost. For the velocity field, free-slip conditions are imposed at the left and top boundaries, while no-slip is enforced at the bottom wall. At the right domain boundary a nonreflective convective outflow condition, we employ the approach proposed by Shields (1936) and Mantz (1973) and use the computed wall shear stress data to check for the possibility of incipient particle motion and/or resuspension. For this purpose, we have to analyze whether or not the "critical" states are reached in which the particles commence their incipient motion along the bottom wall or are being resuspended. This critical state is defined by means of particle Reynolds number \( Re_p \) and nondimensional wall shear stress \( \theta_w \):

\[
Re_p = \frac{\bar{u}_c d_p}{v}.
\]

\[
\theta_w = \frac{\bar{u}_c^2}{(\bar{\rho}_p - \rho) g d_p}.
\]

Here, \( \bar{u}_c \) is the so-called friction velocity related to wall shear stress \( \tau_w \) by

\[
\bar{u}_c = \sqrt{\frac{\tau_w}{\rho}}.
\]

We define the nondimensional friction velocity \( u_t \) using the buoyancy velocity (see Eq. (7)) which is related to the non-dimensional flow quantities by

\[
u_t = \frac{\bar{u}_t}{u} = \left( \frac{C_v \bar{u}}{Re} \right)_{\bar{y}}.
\]

For the given case (C: 0%–F: 100%), the time evolution of the maximum \( u_t \) over the entire bottom wall is shown in Fig. 8. The maximum \( u_t \) observed in the simulation is recorded approximately at time \( t = 4.6 \). This time corresponds to the point when the upper, light current has been reflected from the left wall (Fig. 9). Fig. 9 displays the spatial distribution of the friction velocity over the bottom wall and the corresponding concentration field at \( t = 4.6 \). The maximum friction velocity occurs in the region where the intense vortex in the wake of the current impinges on the bottom wall. The friction velocity has a secondary maximum close to the nose region. To perform our investigation, we employ the maximum value of \( u_t = 0.13 \) and obtain \( \theta_w \approx 0.5 \) and \( Re_p \approx 0.4 \) (see Eqs. (19), (20), and (22)). For any given value of \( Re_p \), Mantz (1973) suggests two thresholds for \( \theta_w \), which are associated with (a) the incipient motion of the particles in the

Fig. 8: Temporal evolution of the maximum friction velocity \( u_t \) computed over the bottom wall. Results are from a two-dimensional simulation of the case (C: 0%–F: 100%) for a bidisperse turbidity current produced by a lock-exchange configuration (same as the setup used in experiments by Gladstone et al., 1998).
form of bed load transport and (b) incipient resuspension of particles, respectively. According to Mantz (1973), at the particle Reynolds number Reₚ = 0.4, the critical value for the onset of resuspension is θₑ ≈ 0.4, which is below the computed value in current simulation. This suggests that particle resuspension may have some influence on the deposit profiles. However, we point out that according to Fig. 8, the maximum friction velocity occurs only in a very short time interval during the early stages of the current’s development. It rapidly drops to values significantly below the critical value of θₑ ≈ 0.4, which suggests the absence of particle resuspension. Second, the deposit layer potentially influenced by the resuspension of particles is limited only to a small spatial region (see Fig. 9). Finally, by time t = 4.6, only 3% of the initial suspended mass is deposited onto the bottom bed, which indicates that only a very small number of particles can potentially be resuspended into the current. On the other hand, the obtained value of θₑ ≈ 0.5 is significantly above the second threshold (θₑ ≈ 0.1) suggested by Mantz (1973) at which the incipient motion of particles on the bottom wall may occur. This observation suggests the possibility of bed load transport of fine particles, which could possibly be responsible for the “shift” observed in the deposit profiles of the last two experimental cases of experiments by Gladstone et al. (1998) (see Fig. 7). In this discussion, however, we have to keep in mind that our computational Reynolds number was lower than the experimental one, which means that the experimental wall shear stress values probably did not reach the computational values. Hence, a definitive conclusion as to whether bed load transport and/or erosion may have occurred in the experiment is not possible. Finally, as we had noted above, there are considerable uncertainties regarding the precise value of the settling velocity, which may also account for some of the observed differences.

In summary, the present depth-resolved simulations confirm the observations of earlier authors, in that they show that even moderate fractions of fine particles cause the current, and with it the coarse particles, to travel significantly further downstream. Fig. 10 shows the runout length Lₜ both for the overall deposit, and for the coarse particles only, for different initial fractions of fine particles. To obtain Lₜ, we search for the location at which the normalized final deposit height Dₚ of the concentration field under consideration (coarse or total) is equal to a prescribed tolerance, e.g., 0.005 in this particular example.

4.3. Polydisperse particle-driven current traveling over a complex bottom surface

As a next step, we focus on a comparison of two-dimensional TURBINS simulation results with two sets of experiments conducted by Kubo (2004). This author conducted a series of experiments to investigate the degree to which turbidity currents are affected by bottom topography, modeled by ramps with slopes similar to those found in deepwater regions. Kubo (2004) also employed the depth-averaged Navier–Stokes equations as a simplified model for such currents.

The first comparison addresses run A4 in Kubo (2004). The experimental suspension is made by mixing siliceous, noncohesive, sand- to silt-sized particles of density ρₛ = 2.65 g/cm². The mixture in the lock is represented by six size fractions with respective diameters of 125.0, 105.1, 88.4, 62.5, 44.20, and 31.25 μm (see Fig. 1 in Kubo, 2004). The relative mass fractions of these particle sizes in the lock are 0.1, 0.15, 0.25, 0.3, 0.15, and 0.05, respectively. The mixture has an overall particle volume fraction of φ_v = 0.02. Kubo (2004) employed the relationship proposed by Gibbs et al. (1971) to relate the settling velocity to the particle diameter. In this way, he obtained settling velocities of 1.136, 0.846, 0.623, 0.33, 0.17, and 0.064 cm/s, respectively. The lock has dimensions L × H × W = 50 × 20 × 17 cm. Using Eqs. (7) and (8), the corresponding buoyancy velocity and Reynolds number are computed as u_b = 18 cm/s and Re = 18,000,
respectively. Our simulations employ a reduced value of $Re = 5,000$. The nondimensional settling velocities are given by 0.0631, 0.047, 0.0346, 0.0183, 0.0094, and 0.0048. The numerical setup is as follows: the domain has a size of $L_x \times L_y = 55 \times 2$, with uniform grid spacing ($\Delta x = 0.039$ and $\Delta y = 0.0125$) in the $x$- and $y$-directions. The velocity and concentration boundary conditions are the same as in the previous section, except that we impose a no-slip condition on the velocity field at the left wall.

Fig. 11 compares the final computational deposit profile obtained with the experimental results for run A4 in Kubo (2004). While the overall trends of the profiles are similar in nature, there exist significant quantitative differences. Specifically, the experimental data show a much longer runout of the current. This is puzzling, as our simulation results had shown much better agreement with the very similar experiments of Necker et al. (2002). Among the potential experimental reasons for the discrepancy could be (a) smaller effective settling velocities than those obtained above from the relationship by Gibbs et al. (1971), (b) the presence of bed load transport, and (c) the presence of erosion. We also note that in Fig. 11 the total areas below the experimental and computational profiles are different. Specifically, the integral over the experimental profile is much smaller than that over the computational profile. Since the figure shows only the section of the flow field downstream of the gate, this indicates that in the experiments substantial deposition occurred already within the lock, which may substantially affect the dynamics of the current downstream of the lock. We will further discuss these issues below.

The second set of experiments conducted by Kubo (2004) includes a nontrivial bottom topography. Fig. 12 shows the schematic of run C5. The geometry includes an initial ramp of size length $l_r = 10$ and height $h_r = 1$ (all dimensions are scaled with the reference length $L_r = 10$ cm). Three successive, identical humps are located downstream of the ramp, with height $s = 0.36$ and length $d = 10$, respectively. As pointed out by Kubo (2004), the resulting slope of 0.036 is of the same order as that of sediment waves produced by turbidity currents in deep-sea regions; cf. Wynn et al. (2000b).

The particle size distribution and settling velocities are identical to those for run A4. In the absence of erosion, the condition imposed at the bottom boundary now reads

$$\mathbf{n} \cdot \nabla c_i = 0, \quad i = 1, \ldots, 6.$$  \hspace{1cm} (23)

The numerical setup and the lock dimensions are identical to those for run A4, except for the length in the $y$-direction which is given by $L_y = 3$.

The final deposit profile obtained from a simulation of this case is compared against experimental run C5 by Kubo (2004) in Fig. 13. Experiments and simulations alike indicate that deposition seems to be stronger in those regions where the flow decelerates due to an upward slope of the bottom topography.

Fig. 14 provides the deposit profiles for the individual particle sizes. We note that the finer particles are least influenced by the given geometry, and therefore can travel past three humps. Similarly to run A4, this case shows significant discrepancies between current simulation results and the experiments by Kubo (2004). As discussed above, factors such as erosion, bed load transport, the lower Reynolds number of the simulation, and uncertainty in the settling velocities may contribute to these differences. Concerning the resuspension of particles, we remark that Kubo (2004) presents theoretical results for both with and without erosion which indicate that erosion does not have a significant effect on the deposit profile (see Fig. 9 in Kubo, 2004).
We note that, given the initial volume fraction of particles ($C_r = 0.02$), the lock dimensions, and the particle density, the total particle mass initially added to the lock is $m_{p0} = 900$ g. On the other hand, if we integrate the area under the curves in Figs. 11 and 13 to obtain the total deposited mass, we obtain for both cases (A4 and C5) a value significantly lower than 900 g, and closer to $m_{p0} = 500$ g. This indicates that a substantial fraction of the particles, close to 45%, settles out before it reaches the lock gate.

The front location as a function of time, shown in Fig. 15, indicates that initially the simulation results duplicate the experimental data very accurately. However, beyond $t = 30$ the computational current slows down significantly as compared to its experimental counterpart. Possible reasons for this discrepancy include the lower Reynolds number in the simulation, the potential presence of erosion, and/or a lower effective settling velocity in the experiment.

Fig. 15 also displays the temporal evolution of the suspended mass for each particle size, normalized by each particle’s initial suspended mass. We note that at the final simulation time of $t = 160$, there are still some particles in suspension, even though the current has come to a nearly complete standstill. When the maximum horizontal particle flux (including advective and diffusive fluxes) in the entire flow field drops below a certain threshold value, we terminate the simulation and “dump” the remaining suspended particles at their current streamwise location. Fig. 16 provides information on the accuracy of this procedure, and on its potential to reduce the computational cost associated with the time integration of the flow field. In this figure, we compare the dumped deposit profiles obtained at two different times $t = 160$ and $t = 300$ with the “most accurate” deposit profile ($D_\infty$) obtained at $t = 1000$ where only 0.2% of the initial particle mass is still in suspension. At $t = 160$, $t = 300$, and $t = 1000$, the maximum recorded horizontal particle fluxes in the entire domain are found to be 0.0047, 0.00067, and 0.000014, respectively. We note that the error due to the dumping of the particles is very small, while this procedure can substantially reduce the computational cost.

4.4. Three-dimensional topography

We now investigate a three-dimensional bidisperse turbidity current flowing past a Gaussian bump in the bottom topography. Our interest focuses on the influence of the bottom topography on the flow structure, and on any resulting nonuniformities of the deposit profile.

Fig. 1 displays the configuration of this problem. The computational domain is given by $L_x / C_{15} = L_y / C_{15} = 17 \times 2 \times 2$, respectively. The center of the Gaussian bump is located at $(x_c, y_c, z_c) = (0.32 L_x, 0.5 L_z)$, while its surface height is given by

$$G(x, z) = h \exp\left(-\frac{(x-x_c)^2+(z-z_c)^2}{2e^2}\right)$$

Here, $h = 0.125 L_y$ denotes the bump height and $e = 0.125 L_y$ controls the width of the bump. The lock has dimensions of...
equal to $u_i$ contains a mixture of two particle sizes, with nondimensional settling velocities respectively.

We set $Re=2000$ and employ uniform grid spacing in all three directions, with $\Delta x = 0.0157$ and $\Delta y = \Delta z = 0.0148$.

$\n \cdot \nabla c_i = 0, \quad i = 1, 2. \quad (25)$

$L_x \times H \times W = 1 \times 2 \times 2$. The suspension includes two particle sizes with nondimensional settling velocities of $u^f_s = 0.035$ and $u^c_s = 0.0035$, and initial relative mass fractions of 90% and 10%, respectively. At the bottom, top, and left boundaries we enforce no-slip, while a free-slip boundary condition is imposed along the two lateral spanwise walls. At the outflow boundary ($x = L_x$), we employ the convective condition discussed earlier (see Eq. (18)) for all flow quantities. For the concentration fields, we enforce a vanishing normal derivative everywhere on the boundaries except for the top wall (see Eq. (16)) and the end wall at $x = L_x$ (see Eq. (18)):

Fig. 17 shows the flow evolution at different times. Initially, it is dominated by the formation of two spanwise startup vortices, cf. the discussion by Härtel et al. (2000b) and Necker et al. (2002). Due to the absence of any initial three-dimensional perturbations, the current maintains its two-dimensional structure until it encounters the bump, at which point the interaction between the front and the bottom boundaries results in the development of strong three-dimensional effects. Once the current has passed the bump, spanwise deformations along the front, known as lobe- and-cleft instabilities (Härtel et al., 2000a), are clearly visible. A strongly three-dimensional flow field, increased turbulent mixing and nonuniformities in sediment deposition (Necker et al., 2005) are direct consequences of these instabilities on the flow.

Clearly, nonuniform effects including depletive/accumulative flow regimes mostly arise where the current experiences constrictions and/or expansions. Kneller and Branney (1995) showed that such effects can have an influence on the spatial distribution of turbidites. The topography of the Gaussian bump causes the current to be deflected laterally, thus resulting in nonuniformities of the deposit profile (see Fig. 18). The figure shows that these nonuniformities are more pronounced for the coarse than for the fine particles. In this context we remark that Al Ja Aidi (2000) showed that if the obstacle height is more than 10% of the current’s height, the flow develops a dividing streamline. The flow sections below this dividing streamline are deflected to the sides of the obstacle, while the upper portions of the flow pass over the obstacle (see Fig. 17).

Fig. 18 compares the streamwise deposit profiles in two planes, the midplane ($z = 0.5L_z$) and at the edge of the bump ($z = 0.15L_z$), with the case of a flat surface. While the deposition of particles near the bump decreases strongly in the midplane, it increases closer to the lateral walls. This is consistent with the lateral motion of the current. Fig. 19b investigates the influence separately for fine and coarse particles. It indicates that the influence of the bottom topography is more pronounced for the larger particles.

Fig. 20 displays the spanwise deposit profiles of the coarse particles at different $x$-locations. It confirms the significant influence of the topography on the deposit profiles near the bump, in terms of reducing deposition in the midplane while enhancing it near the spanwise boundaries. This effect progressively diminishes in the far wake.

5. Summary and conclusions

In an earlier publication (Nasr-Azadani and Meiburg, 2011) we had demonstrated the accuracy of the software tool TURBINS for simulating constant density and gravity-driven fluid flows in complex geometries. Here, we have discussed the extension of TURBINS for simulating particle-driven flows. Comparisons of such quantities as deposit profiles, front location, and runout...
length with the experimental and computational data of de Rooij and Dalziel (2001) and Necker et al. (2002) demonstrate the ability of TURBINS to accurately reproduce corresponding experiments. We extended the investigation of Necker et al. (2002) by focusing on the influence of the conditions imposed at the left side wall and the top boundary.

We subsequently focused on bidisperse turbidity currents, by comparing with the experiments of Gladstone et al. (1998). Varying the initial mass fractions of coarse and fine particles, we observe that the runout length and final deposit profiles of the coarse particles are altered significantly by the influence of the fine particles. This clearly indicates the coupling between the different particle sizes in polydisperse mixtures and their effect on the deposit profiles. We also carried out a detailed analysis using the maximum friction velocity to assess the occurrence of resuspension of particles and also bed load transport. We employed the approach by Mantz (1973) and found that the discrepancies appearing in the final deposit profiles are more likely due to bed load transport than to resuspension of particles.

In order to investigate the effect of complex bottom topography on turbidity currents, we compared with the experiments of Kubo (2004). While we find qualitative agreement, there are quantitative discrepancies. Among the reasons for these may be an overestimation of the experimental particle settling velocity (cf. also the comments in Gladstone et al., 1998), along with the presence of bed load transport and/or erosion.

To demonstrate the three-dimensional capabilities of TURBINS, we study the deposit profiles obtained for turbidity currents flowing over a Gaussian bump. As the current interacts with the bump, we observe the generation of so-called lobe-and-cleft instabilities. Furthermore, the flow is deflected laterally, which results in the development of nonuniformities in the deposit profiles especially in the near wake region of the bump. Deposit profiles plotted in various spanwise and streamwise planes suggest that the fine particles are less influenced by the obstacle geometry than the coarse particles.

In summary, TURBINS has been demonstrated to be a versatile numerical tool for investigating the interactions between turbidity currents and seafloor topography. The implementation of a turbulence model in TURBINS is currently underway.

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References


Fig. 19. Influence of the Gaussian bump on the deposit profile of a bidisperse turbidity current. (a) Comparison of the total final deposit profiles against a similar current flowing over a flat surface. (b) Influence of the bump on the deposition of coarse and fine particles, respectively. All deposit heights \( D_\alpha \) are normalized by the corresponding total initial mass multiplied by \( W \). Note that solid lines represent the edge of the bump \((z = 0.15L_z)\) and dashed lines correspond to the midplane \((z = 0.5L_z)\), respectively. The Gaussian bump causes significant spanwise nonuniformities in the deposition profile of the current propagating over the bump.

Fig. 20. Spatial variation of the final deposit profile for the coarse particles, for a bidisperse turbidity current flowing past a Gaussian bump. For each streamwise location, the deposit profile is plotted along the spanwise direction \( z \). Curves (1–6) correspond to \( x = 4.8, 5.5, 6.2, 6.9, 7.6, \) and \( 8.3 \), respectively. The dashed line represents the cross section of the bump at \( x = 5.5 \). Note that all \( D_\alpha \)'s are normalized by total initial mass of coarse particles multiplied by \( L_s \).


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