Buoyant particles and fluid turbulence

Varghese Mathai

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BUOYANT PARTICLES AND FLUID TURBULENCE

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1

Introduction

1.1 Fluid turbulence

Turbulence is a common phenomenon in everyday life - we notice it all around us, in hurricanes and in waterfalls, in the breaking waves in the oceans, and even during the typical flight journey (see examples in Fig. 1.1). Many scientists and engineers have devoted a lifetime of work in pursuit of the subject. Yet, it remains one of the least understood phenomena in classical physics.

When addressing turbulence, it is typical to describe the features that characterize it. Turbulence in fluids may be regarded as irregular, random, and chaotic. It is also characterized by the presence of a range of length- and time-scales, which enables it to transport and mix the fluid very effectively. In addition, turbulence is dissipative by nature and requires energy input to sustain, in the absence of which it simply decays. This also makes it time-irreversible, meaning that, statistically, one can tell the arrow of time from a turbulent flow signal [1, 2].

Since the first pipe flow experiments by Osborne Reynolds in 1883, turbulence has been quantified using the Reynolds number $Re \equiv U D/\nu$, where U is a fluid velocity scale, D is the pipe diameter, and ν is the kinematic viscosity of the fluid. This dimensionless number presents the ratio between inertial and viscous forces in the flow under consideration. When $Re \gg 1$, the flow is in general regarded as turbulent. Subsequently, several different definitions of Re have been proposed, ranging from the Taylor-Reynolds number Re_λ for homogeneous turbulence, to the friction Reynolds number Re_τ for pipe and channel

Figure 1.1: Examples of turbulent flows with different scales of motion. (a) A breaking wave in the ocean. Satellite images of (b) ocean currents (released by NASA), and (c) Hurricane Katrina. (d) Schlieren image of human thermal plume [3].

flows. While quantitatively different, they may be regarded equivalent, since they can be linked to the original definition.

From a mathematical point of view, fluid turbulence is described by the *Navier-Stokes equations* [4]. However, the non-linearity of the equations renders them difficult to treat. Therefore, one of the alternatives was to resort to a statistical description of turbulence. This has proven to be one of the most successful approaches in the last century at describing turbulent flows in general. On this front, many important ideas have been proposed. Among these, the vision of Andrei Kolmogorov has been one of the most influential. Kolmogorov's theory introduced the hypothesis of a universal, isotropic, homogeneous statistical distribution for the small scales of motion in high-Reynolds-number turbulence [5]. The theory proposed the idea that a turbulent flow is composed of *eddies* of different length and time scales, with a cascade of energy from the large to the small scales driving the turbulence. The $-5/3$ law of the energy spectrum, which emerges out of this hypothesis, has been found in a range of flows, including boundary layers and shear flows, and even in situations where there are substantial departures from isotropy [6]. This powerful vision, to this day, influences our understanding of fluid turbulence.

1.2 A particle in a turbulent flow

While Kolmogorov's vision of a universal turbulent flow might be elegant in concept, many real flows are subject to additional complexities. A wide range of flows in nature and industry are turbulent and at the same time contain dispersed particles. Typical examples are pollutants dispersed in the atmosphere, bubbles and planktons in the oceans, sediment-laden river flows, particulate eruptions from volcanos, and fuel sprays in combustion engines (see some examples in Fig. 1.2). Many interesting physical phenomena occur in such particle-laden flows. For example, in the oceans, the blooming of plankton results in interesting patterns, where the plankton cluster or distribute unevenly (see Fig. 1.2(c)) despite the carrier flows being nearly homogeneous at the corresponding length scales. In other situations, buoyant particles and bubbles rise through oceans, creating localized mixing and concentrations of dissolved gases (see Fig. 1.2(b)). The presence of these have been correlated with measurements of warmer waters and larger temperature gradients in certain parts of the oceans [7, 8]. In addition, bubbles and particles are commonly used to induce liquid agitation in multiphase flows. This finds application in the chemical industry, where bubble-column reactors enhance the mixing and heat transfer through the liquid. A simplified picture of single-phase turbulent flow is not adequate to model and understand these flows. This makes 'particles in turbulence' an active area of research relevant to a diverse number of fields.

The interactions between particles and turbulent flows can range from being 'simple', in some situations, to exceedingly 'complex' in others. In many practical environments, particles are small and their suspensions may be regarded as dilute (volume fraction $\phi_v < 0.1\%$). In these cases, often a simplified problem is considered, that of a rigid spherical particle in a homogeneous isotropic turbulent flow. A proper description of the system requires knowledge of several parameters, which define the flow as well as the particles. The state of homogeneous turbulence is fully determined by the kinematic viscosity, and by a length and a time scale of the flow. With the introduction of particles, we bring in the particle diameter d_p and the particle density ρ_p , the presence of which makes the fluid density ρ_f also important. By dimensional considerations, one can form three dimensionless groups: $\Xi \equiv d/\eta$, $\Gamma \equiv \rho_p/\rho_f$, and Re_λ ; the specification of these is expected to fully define the problem*. Several studies have been performed in this domain of dilute suspensions of passively advected particles in turbulence [9]. Within this framework, the behavior of

*Note that the Stokes number St is also an important parameter. St gives a measure of how slowly the particle responds to the flow. However, with Ξ and Γ specified, St is fixed.

Figure 1.2: (a) A smog-hit city. (b) CO_2 bubbles rising from vents in the ocean floor. (c) A plankton bloom in the ocean. (d) Fuel spray emanating from a generic aircraft engine fuel injector. Picture Credits: (a), (b) David Liittschwager/National Geographic Creative, and (c) MERIS: plankton blooms: ESA, and (d) GE reports.

the particle is often compared and contrasted with the movement of a fluid element, and inertial range scaling arguments based on Kolmogorov's theories have been widely used [10–12].

The above description of a dilute suspension of passively advected particles, although widely in use [9, 13], is valid only when gravity is negligible. When gravity is not negligible, the Froude number $\text{Fr} \equiv a_\eta/g$ [14] becomes important. This measures the relative strength of the turbulent acceleration a_η in the flow to the acceleration due to gravity g . When Fr is small, the effects of gravitational settling become important, which induces a range of new physical phenomena, including clustering, enhanced settling speeds, and reduced dispersion for settling particles in turbulence [14–16].

When the particle size is large, often the approximation of a passively advected particle is no longer valid. This would be the case when the particle Reynolds number Re_p is finite. This can occur in two situations. First,

when the turbulent velocity fluctuations u_{rms} at the scale of the particle are large enough to result in finite $\text{Re}_p \equiv u_{rms}d_p/\nu$. Alternately, gravitational rise/settling of a particle could lead to a finite $\text{Re}_p \equiv V_r d_p/\nu$, where V_r is the relative rise/fall velocity of the particle. Many problems of practical interest fall within the above categories. In these situations, the turbulent eddies influence the particle motion, and the particle back-reacts on the flow, often leading to major modifications of the particle and the flow dynamics. For example, bubbles and/or particles can be used to modify the drag and the transport of heat in turbulent flows [17, 18]. Similarly, the presence of buoyant bubbles/particles can affect the cascade of energy in turbulent flows [19]. Few investigations have addressed the dynamics of such strongly interacting buoyant particles carried by turbulent flows.

1.3 Turbulent flow experiments

The experiments in this thesis were mostly performed in the Twente Water Tunnel facility (TWT). A drawing of the facility is shown in Fig. 1.3. The setup houses an active grid, which generates nearly homogeneous and isotropic turbulence in the measurement section downstream [20–23]. The water tunnel is symmetrical about the horizontal plane passing through the mid-height of the measurement section. Therefore, it was possible to have two equivalent flow arrangements: (a) a downward flow with the active grid on top, and (b) an upward flow with the active grid placed below the measurement section. As you will see in the chapters of this thesis, we switch between these two arrangements in order to gain certain experimental advantages.

1.4 A guide through this thesis

In this thesis, we address how *buoyancy* affects the dynamics and wakes of particles and bubbles in a range of flow environments. The primary focus is on particle- and bubble-laden turbulent flows. We look into the motion of light[†] and heavy particles, ranging from microbubbles to finite-sized rigid light spheres, as well as large deformable air-bubbles, all of which experience a net buoyancy force when submerged in the fluid. For small particles and bubbles, the dynamics is essentially one-way coupled. However, for millimetric bubbles and finite-sized spheres, the particle wakes influence their dynamics.

In Chapter 2, we present the acceleration statistics of *tiny* bubbles dispersed in a turbulent water flow. We find that bubbles have a higher acceleration vari-

[†]In this thesis, the terms *light* and *buoyant* are used interchangeably to refer to particles that are less heavy than the carrier fluid.

Figure 1.3: The Twente Water Tunnel, which is the main facility used for the experiments reported in the present thesis. The flow direction and the location of the active grid were switched for convenience of the study.

ance compared to tracer particles, which occurs despite their small size and minute Stokes numbers. Using the assumption of *one-way coupling*, we theoretically and numerically show that the enhanced accelerations could originate due to the vertical drift of these particles through the turbulent vortices in the flow. In addition, we predict the changes in the acceleration statistics (variance, time-correlation, and intermittency) for arbitrary density particles in homogeneous and isotropic turbulence.

In Chapter 3, we move our attention to the case of *finite-sized* buoyant spheres in turbulence. In this case, in addition to the accelerations caused by the vertical drift, a stronger contribution to acceleration originates from the wake-induced forces on the particles. We show that in this regime, point-particle models fail to predict the acceleration behavior even after finite-size corrections [24] are included.

In Chapter 4, we consider a *finite-sized* buoyant sphere, but in addition to

its translation, we track its instantaneous rotational motions. To this end, we develop an efficient experimental method to track simultaneously the translation and rotation of a spherical particle in turbulence. Results are presented on the accuracy and numerical scalability of the method, in addition to experiments in the Twente Water Tunnel.

In Chapter 5, we employ the technique developed in Chapter 4 to study the effect of rotation on the translational dynamics of finite-sized rigid light spheres. We find that a change in the rotational freedom of the particle can be effected by changing its rotational inertia. This can produce significant changes to the particle's translational dynamics. Thus, the moment of inertia is shown to be an important control parameter in problems concerning the translational dynamics of buoyant particles in both turbulent and quiescent flow environments.

In Chapter 6, motivated by observations of moment of inertia induced effects for rising spheres (in Chapter 5), we numerically study the oscillatory dynamics and wakes of freely rising and falling cylinders. We establish two regimes of motion, demarcated by a transitional Strouhal number St_t . For $St < St_t$, the mass of the particle determines the oscillation amplitude, while for $St > St_t$, the dynamics is governed primarily by the rotational motion of the particle, which manifests through the particle moment of inertia[‡].

In Chapter 7, we study the damped oscillations of a cylindrical pendulum in water. We model the dynamics by accounting for the pendulum mass, buoyancy, added mass, and bearing friction to predict the oscillatory amplitudes and frequencies for a wide range of mass-density ratios. The predictions are in reasonable agreement with our experimental observations for a range of mass-density ratios.

In Chapter 8, we move from rigid particles to the case of deformable air-bubbles dispersed in a turbulent water flow. We study how the buoyancy-induced motion of these bubbles influences the agitation in the liquid phase. We show that the *bubble* parameter [19] can fairly capture the increase in the liquid agitation.

Finally, we conclude and summarize the work done in this thesis.

[‡]Note that St refers to the Stokes number, while St refers to the Strouhal number [25]. These appear in some of the chapters of this thesis and should not be mistaken for each other.

2

Gravity-induced accelerations of microbubbles in turbulence *

We report on the Lagrangian statistics of acceleration of small (sub-Kolmogorov) bubbles and tracer particles with Stokes number $St \ll 1$ in a turbulent flow. At decreasing Reynolds number, the bubble accelerations show deviations from that of tracer particles, i.e. they show a larger acceleration variance and a quicker acceleration decorrelation despite their small size and minute St . Using direct numerical simulations, we show that these effects arise due the drift of these particles through the turbulent flow. We theoretically predict this gravity-induced effect for developed isotropic turbulence, with the ratio of Stokes to Froude number or equivalently the particle drift-velocity governing the enhancement of acceleration variance, and the reductions in correlation time and intermittency. The present findings are relevant to a range of scenarios encompassing tiny bubbles and droplets that drift through the turbulent oceans and the atmosphere.

*Based on: V. Mathai, E. Calzavarini, J. Brons, C. Sun, and D. Lohse, Microbubbles and microparticles are not faithful tracers of turbulent acceleration, Phys. Rev. Lett. 117, 024510 (2016).

Experiments by Mathai and Brons. Numerical simulations by Calzavarini. Analysis and writing by Mathai and Calzavarini. Supervision by Sun and Lohse, discussion of results and proofreading by everyone.

2.1 Introduction

Heavy and light particles caught up in turbulent flows often behave differently from fluid tracers. The reason for this is usually the particle's inertia, which can drive them along trajectories that differ from those of the surrounding fluid elements [9, 26–28]. Due to their inertia, measured by the Stokes number $St^\dagger$, such particles depart from fluid streamlines and distribute nonhomogeneously even when the carrier flow is statistically homogeneous [9, 12, 29–34]). Numerical studies have captured several interesting effects of particle inertia through point-particle simulations in homogeneous isotropic turbulence [12, 35–37]. For instance, with increasing inertia, light particles showed an initial increase in acceleration variance (up to a value $\langle a^2 \rangle \sim 9$ times the tracer value) followed by a decrease, while heavy particles showed a monotonic trend of decreasing acceleration variance [24]. Such modifications of acceleration statistics arose primarily from the slow temporal response of these inertial particles, i.e. when St was finite [32, 38–41]. In contrast, a lower limit of inertia can be imagined ($St \ll 1$), when the particles respond to even the quickest flow fluctuations and, hence, are often deemed good trackers of the turbulent flow regardless of their density ratio [9, 24, 38, 42]. The widespread use of small bubbles and droplets in flow visualization and particle tracking setups (e.g. Hydrogen bubbles in liquid flows, and oil droplets in gas flows) is founded on this one assumption – that $St \ll 1$ renders a particle responsive to the fastest fluctuations of the flow [21, 43–46].

In many practical situations, particles are subjected to body forces, typically gravitational, and in some situations centrifugal [17, 47]. This can be the case for rain droplets and aerosols settling through clouds, and for tiny air bubbles and plankton drifting through the oceans [48, 49]. Due to their density mismatch, these particles either settle or rise in the flow. The effects of gravitational settling were first brought to light through numerical studies of particles in random and cellular flow fields [15, 16, 50–54]. These showed preferential concentration and enhanced settling velocities for heavy particles settling in random and cellular flow fields. More recently, inertial effects on settling particles were analysed using direct numerical simulations of fully developed homogeneous isotropic turbulence [14, 39, 40, 55, 56]. These revealed that gravity can lead to major modifications of particle clustering, relative velocity and pair statistics, which could be characterized as a function of St and the ratio of turbulent to gravitational acceleration: a_η/g .

While the effects of gravity on particle settling velocity and clustering were

[†]Stokes number $St \equiv \tau_p/\tau_\eta$, where τ_p is the particle response time, and τ_η is the Kolmogorov time scale of the flow.

shown to be significant [14, 15], another crucial observable of Lagrangian turbulence is the particle's acceleration statistics (eg. variance, correlation and intermittency). Acceleration is important because its variance and time correlation may be linked to the energy dissipation rate, a quantity central to characterizing turbulent flows. Yet another feature unique to turbulent flows is the high level of intermittency, or deviation from Gaussian statistics. Acceleration flatness is key to quantifying this for turbulent flows. Therefore, a generic description of these quantities for rising and settling particles of arbitrary density is desirable.

In this Chapter, we present the Lagrangian acceleration statistics of small air-bubbles and neutrally buoyant tracer particles in a turbulent water flow where the Taylor-Reynolds number Re_λ is varied in the range 130 – 300. With decreasing Re_λ the bubble accelerations show deviations from that of tracer particles, which occur despite their very small St (0.004 – 0.017) and small particle size. To explain these, we conduct DNS of particles in homogeneous isotropic turbulence, which reveals that the deviations could arise due to the drift of these particles through the turbulent flow. We make generic predictions of the gravity-induced deviations for an arbitrary density particle, with the ratio St/Fr or equivalently the ratio of particle drift velocity to Kolmogorov velocity, as the relevant parameter controlling the deviations from ideal tracer behavior. Further, we provide insight into the modification of intermittency of particle acceleration arising due to gravity.

2.2 Experiments

The experiments were performed in the Twente Water Tunnel facility (TWT), in which an active grid generated nearly homogeneous and isotropic turbulence in the measurement section [22]. The turbulent flow was characterized using hot-film anemometry technique at different Re_λ (see Table 2.1). Small bubbles ($\approx 150 \pm 25 \mu\text{m}$) were generated by blowing pressurized air through a porous ceramic plate. The particles were imaged using a high-speed camera (Photron PCI-1024) at a recording rate of 1000 fps. The camera moved on a traverse system, and illumination was provided using a 50 W Pulsed Laser (Litron LDY-303HE) (see Fig. 2.1). Similar moving camera setups have been used with heavy particles in a wind tunnel [57], however, the tracking duration was short in those experiments. In the present case, we introduce an experimental modification, which allows for long particle trajectories to be recorded. We placed a mirror inside the tunnel at 45° inclination with the horizontal (see Fig. 2.1). The laser beam was expanded into a volume, which was reflected vertically by the mirror. Recording was stopped when the cam-

Figure 2.1: Schematic of the measurement section of the Twente Water Tunnel. Bubbles and neutrally buoyant tracer particles of diameters $\approx 150 \pm 25 \mu\text{m}$ and $\approx 125 \mu\text{m}$, respectively, were dispersed in the flow for particle tracking experiments.

era moved closer than 0.4 m from mirror. At this distance, flow disturbance due to the downstream mirror was expected to be minor. The particles were tracked to obtain their position versus time from the experiments. As discussed earlier, the focus of the present study is on the acceleration statistics. Obtaining acceleration from position data requires accurate determination of higher derivatives. Here we have combined the Gaussian-kernel smoothing method [58] with a smoothing-spline based method [59] which eliminated biases due to a priori choice of filter windows and ensured reliable estimates of the acceleration.

2.3 Results

We first address the question of how the bubble accelerations compare to those of similar-sized neutrally buoyant particles at different Re_λ . According to the prediction by Heisenberg and Yaglom [11], the single-component variance of acceleration should follow the relation $\langle a^2 \rangle = a_0 \epsilon^{3/2} \nu^{-1/2}$, where ϵ is the dissipation rate, and ν the kinematic viscosity. In Fig 2.2(a), we plot the Heisenberg-Yaglom constant a_0 for bubbles along with that for tracer particles. At high Re_λ , the bubbles behave similarly to tracers, with comparable a_0 . However, at low Re_λ the bubbles show deviations from tracers, with an elevated a_0 . The horizontal acceleration shows the greatest deviation, with $a_0 \approx 6$, while for the vertical component, $a_0 \approx 4.5$ at $\text{Re}_\lambda \approx 134$. For neutrally buoyant tracer particles, a_0 is lower, ≈ 2.1 at $\text{Re}_\lambda \approx 134$, and shows a marginal increase with Re_λ . Thus, the horizontal component of the acceleration variance for bubbles is almost three times that of the tracer value at the lowest Re_λ (see Fig. 2.2(b)). This is also reflected in the correlation time for

Figure 2.2: (a) Heisenberg-Yaglom constant, a_0 , estimated for bubbles and tracer particles from experiments at different Re_λ . The dashed curve gives the a_0 estimate according to [60]. (b) Normalized acceleration variance and (c) correlation time of acceleration for bubbles vs Re_λ from experiments. $\langle a_T^2 \rangle$ is the tracer particle acceleration variance. τ_p and τ_T are defined as the 0.5 crossing time of correlation for bubbles and tracer particles, respectively. Inset: Normalized correlation function of acceleration, $C_a(\tau)$, at the lowest $Re_\lambda \approx 134$.

bubbles (Fig. 2.2(c) & Inset), which is shorter compared to that of tracers. Thus, the mismatch between acceleration statistics of bubbles and neutrally buoyant tracer particles increases with decreasing Reynolds number. However, the lowest Re_λ corresponds also to the smallest St in our experiments (see Table. 2.1) Therefore, these deviations could not possibly be due to the particle's inertia [9].

Table 2.1: Flow characteristics in the Twente Water Tunnel. Re_λ – the Taylor-Reynolds number (approximate), St – the bubble Stokes number, and a_η/g the ratio of turbulent to gravitational acceleration [14].

Re_λ	134	153	188	263	301
St	0.004	0.005	0.008	0.011	0.017
a_η/g	0.0016	0.0026	0.0044	0.0072	0.0141
St/Fr	4.899	4.159	3.498	2.953	2.363

2

2.4 Particle equation of motion

From Fig 2.2(b), we also note that the vertical component of acceleration is consistently lower as compared to the horizontal one. This anisotropy is not inherent in the carrier flow [20] and therefore suggests the role of gravity [52]. We note that with decreasing Re_λ , the ratio of turbulent to gravitational acceleration, a_η/g , decreases in our experiments (Table. 2.1). In order to investigate this effect in a systematic way, we perform DNS of homogeneous isotropic turbulence at $\text{Re}_\lambda \approx 80$ in the presence of gravity. For the particles, we use a model considering a dilute suspension of passively advected point-spheres acted upon by inertial and viscous (Stokes drag) forces. The model equation of motion for a small inertial spherical particle advected by a fluid flow field, with velocity $\mathbf{U}(\mathbf{X}(\mathcal{T}), \mathcal{T})$, is:

$$\mathcal{V} \rho_p \ddot{\mathbf{X}} = \mathcal{V} \rho_f \frac{D\mathbf{U}}{D\mathcal{T}} + \mathbf{F}_{\text{AM}} + \mathbf{F}_{\text{D}} + \mathbf{F}_{\text{B}} \quad (2.1)$$

where $\mathcal{V} = \frac{4}{3}\pi a^3$ is the particle volume, with a being the particle radius, and ρ_f and ρ_p the fluid and particle mass densities, respectively. The forces contributing on the right-hand-side besides the fluid acceleration (which includes the pressure gradient term) are the added mass $\mathbf{F}_{\text{AM}}$, the drag force $\mathbf{F}_{\text{D}}$, and the buoyancy $\mathbf{F}_{\text{B}}$ [61–63]:

$$\mathbf{F}_{\text{AM}} = \mathcal{V} \rho_f C_M \left(\frac{D\mathbf{U}}{D\mathcal{T}} - \ddot{\mathbf{X}} \right), \quad (2.2)$$

$$\mathbf{F}_{\text{D}} = 6 \pi \mu a (\mathbf{U} - \dot{\mathbf{X}}), \quad (2.3)$$

$$\mathbf{F}_{\text{B}} = \mathcal{V} (\rho_f - \rho_p) g \hat{\mathbf{e}}_z, \quad (2.4)$$

where μ is the dynamic viscosity, g is the gravity intensity, and $\hat{\mathbf{e}}_z$ is the unit-vector in the vertical direction. Note that we use the ideal flow (inviscid) added mass coefficient for a sphere, i.e $C_M = 1/2$. This leads to

$$\ddot{\mathbf{X}} = \frac{3\rho_f}{\rho_f + 2\rho_p} \left(\frac{D\mathbf{U}}{D\mathcal{T}} + \frac{12\nu}{d_p^2} (\mathbf{U} - \dot{\mathbf{X}}) + g \hat{\mathbf{e}}_z \right) - g \hat{\mathbf{e}}_z, \quad (2.5)$$

Figure 2.3: Normalized acceleration variance for buoyant particles vs a_η/g obtained from Eulerian-Lagrangian DNS at $\text{Re}_\lambda \approx 80$. Hollow and solid symbols correspond to horizontal and vertical components, respectively.

where $\nu \equiv \mu/\rho_f$ is the kinematic viscosity, and d_p is the particle diameter. Here we neglect lift, history, and finite-size Faxén forces, since these are verified to be small in point-particle limit and when the particle size is smaller than the Kolmogorov length scale η of the flow [24, 64, 65].

The particles under consideration are buoyant ($0 \leq \Gamma \equiv \rho_p/\rho_f < 1$) and are chosen to have a very small Stokes number[‡] ($\text{St} \approx 0.05$). We vary the gravity intensity g for these particles, resulting in a range of values for a_η/g , the Froude number definition in [14, 40]. We first address the case of bubbles ($\Gamma = 0$ in Fig. 2.3) at various strengths of g . With increasing g , we recover the trends observed in our experiments, i.e. the bubble acceleration variance increases. Gravity enhances the acceleration in both vertical and horizontal directions.

Fig. 2.3 clearly demonstrates the role of gravity in enhancing the particle's acceleration variance. At the same time, we note that the degree of enhancement diminishes with growing Γ even at fixed a_η/g . Also, the effect of gravity on $\langle a_p^2 \rangle / \langle a_T^2 \rangle$ appears more pronounced in our simulations (see Fig. 2.2(b)) and Table 2.1). Thus, the Froude number, if defined as a_η/g [14, 40], can only give a qualitative prediction of the gravity effect.

For a better appreciation of the contribution of gravity to the particle dynamics, we non-dimensionalize eq. (2.5) in terms of the Kolmogorov length η and time scales τ_η . We obtain

$$\ddot{\mathbf{x}} = \beta \frac{D\mathbf{u}}{Dt} + \frac{1}{\text{St}}(\mathbf{u} - \dot{\mathbf{x}}) + \frac{1}{\text{Fr}}\hat{\mathbf{e}}_z, \quad (2.6)$$

[‡]Stokes number can be defined in several ways. Here we have used the expression $\text{St} \equiv d_p^2/(12\beta\nu\tau_\eta)$, where $\beta \equiv 3\rho_f/(\rho_f + 2\rho_p)$ is the effective mass density ratio.

where $St \equiv d_p^2/(12\beta\nu\tau_\eta)$ is the Stokes number and $Fr \equiv a_\eta/((\beta - 1)g)$ is a buoyancy-corrected Froude number that takes the particle density, through $\beta \equiv 3\rho_f/(\rho_f + 2\rho_p)$, into account. In this situation, two important small-Stokes limits may be considered [39, 52]. At high turbulence intensities ($Fr \rightarrow \infty$), the third term on the right-hand-side of eq. (2.6) may be neglected. Under this condition the vanishing St limit leads to $\dot{\mathbf{x}} \simeq \mathbf{u}$ for the velocity. This leads to the well known result $\ddot{\mathbf{x}} \simeq D_t \mathbf{u}$ for particle acceleration, where $D_t \mathbf{u}$ is the fluid tracer acceleration. However, for finite Fr , the small St limit leads to $\dot{\mathbf{x}} \simeq \mathbf{u} + \frac{St}{Fr} \hat{\mathbf{e}}_z$. For the particle acceleration this implies:

$$\ddot{\mathbf{x}} \simeq D_t \mathbf{u} + \frac{St}{Fr} \partial_z \mathbf{u} \quad (2.7)$$

2.4.1 Acceleration variance

From eq. (2.7), one can obtain an approximate form for the single-component acceleration variance. These are

$$\langle \ddot{x}^2 \rangle \simeq \langle (D_t u_x)^2 \rangle + \left(\frac{St}{Fr} \right)^2 \langle (\partial_z u_x)^2 \rangle, \quad (2.8)$$

$$\langle \ddot{z}^2 \rangle \simeq \langle (D_t u_z)^2 \rangle + \left(\frac{St}{Fr} \right)^2 \langle (\partial_z u_z)^2 \rangle \quad (2.9)$$

where x is one of the horizontal components, and z is the vertical component. Note that the linear terms in St/Fr vanish under the assumption of no correlation between terms of the type $\mathbf{u} \cdot \nabla u_i$ and $\partial_z u_i$, i.e. there is no instantaneous correlation between the velocity field and its gradient. Under isotropic turbulent conditions, the following relations are verified [4]:

$$\langle (\partial_z u_x)^2 \rangle \simeq \frac{2}{15} \frac{\epsilon}{\nu}, \quad (2.10)$$

$$\langle (\partial_z u_z)^2 \rangle \simeq \frac{1}{15} \frac{\epsilon}{\nu}, \quad (2.11)$$

$$\langle (D_t u_i)^2 \rangle = a_0 \epsilon^{3/2} \nu^{-1/2}, \quad (2.12)$$

where i denotes one of the components x , y , or z , and a_0 is the so-called Heisenberg-Yaglom constant. From this, one obtains the relations linking the acceleration variance of the particles to that of the fluid tracers:

$$\frac{\langle a_h^2 \rangle}{\langle a_T^2 \rangle} \equiv \frac{\langle \ddot{x}^2 \rangle}{\langle (D_t u_x)^2 \rangle} \simeq 1 + \frac{2}{15 a_0} \left(\frac{St}{Fr} \right)^2, \quad (2.13)$$

$$\frac{\langle a_v^2 \rangle}{\langle a_T^2 \rangle} \equiv \frac{\langle \ddot{z}^2 \rangle}{\langle (D_t u_x)^2 \rangle} \simeq 1 + \frac{1}{15 a_0} \left(\frac{St}{Fr} \right)^2, \quad (2.14)$$

Figure 2.4: Normalized acceleration variance in horizontal (h) and vertical (v) directions vs St/Fr for bubbles from our experiments. The dashed lines give the predictions based on the eq. (2.13) & (2.14), using a_0 obtained for the tracer particles in the present experiments (Fig. 2.2(a)).

where a_h and a_v are the horizontal and vertical accelerations, respectively, for an arbitrary-density particle, a_T is the tracer particle acceleration, and x and z represent the horizontal and vertical directions, respectively.

In Fig. 2.4, we compare the normalized acceleration variance vs St/Fr from experiment with the predictions (eq. (2.13) & (2.14)). The dashed lines show the theoretical predictions using the a_0 from the present experiments (Fig. 2.2(a)). The experimental data-points are in reasonable agreement with our predictions. We can conclude that the apparent Re_λ dependence that was seen in our water tunnel experiments (Fig. 2.2) is in fact a St/Fr effect, since the St/Fr increases with decreasing Re_λ in our experiments (see Table 2.1).

The present bubble-tracking experiments cover a narrow range of St/Fr in the interval (2, 5) at fixed density-ratio ($\beta = 2.99$), while eq. (2.13) & (2.14) should be valid for arbitrary density-ratio. To test this, we compare the results of DNS for an extended range of density-ratios $\beta = [0, 3]$ and $St/Fr = [-10, 20]$. In Fig. 2.5(a), the left half ($St/Fr < 0$) points to heavy particles, and the right one ($St/Fr > 0$), to buoyant particles. We obtain a good agreement between the predictions and the simulations. It is important to note that in both the experiments and the simulations, a_0 varies with Re_λ . However, for sufficiently large Reynolds numbers, one may expect a_0 to be practically constant [11]. In these situations the ratio St/Fr becomes the sole control parameter governing the enhancement of acceleration variance. Therefore, the present results might have broad applicability, to even large Re_λ atmospheric and oceanic flows.

Figure 2.5: (a) Normalized acceleration variance, (b) normalized correlation time, and (c) normalized flatness factor vs St/Fr for a family of buoyant and heavy particles, obtained from DNS. Solid and dashed curves in (a) show our theory-predictions for horizontal and vertical accelerations, respectively (eq. (2.13) & (2.14)). (b) The black curve shows the theoretical prediction for correlation-time. (a)-(c) Hollow and solid symbols correspond to horizontal and vertical components, respectively.

2.4.2 Acceleration time-correlation

In Fig. 2.5(b), we plot the simulation results for the evolution of the Lagrangian time-correlation of acceleration with the ratio St/Fr . The left branch points to heavy particles, and the right one, to light particles. With increasing magnitude of St/Fr , we observe a decline in the correlation time for both heavy and light particles. Clearly, the drifting of the buoyant or heavy particle

Figure 2.6: (a)-(d) Normalized autocorrelation function for fluid tracers and bubbles at different St/Fr values. Here, a_x and a_z are the horizontal and vertical accelerations, respectively. $\partial_z u_x$ and $\partial_z u_z$ are the spatial-velocity gradients, which contribute to the a_x and a_z , respectively. In (b)-(d) St/Fr increases from 4.5 to 20.

through the flow affects the correlation time. We model this by considering a particle drifting through the flow at a speed u_r . In the absence of particle drift the decorrelation time of a neutrally buoyant tracer may be given by τ_T [66]. Based on the characteristic velocity of a particle in the turbulent flow, u_{rms} , we estimate the length scale corresponding to this decorrelation time as $\Lambda \sim u_{rms}\tau_T$. Now, for a buoyant or heavy particle, the time of correlation is reduced due to an extra drift speed. Therefore, the new correlation time may be written as $\tau_p \approx \tau_T/(1 + u_r/u_{rms})$. For homogeneous isotropic turbulence the u_{rms} may be expressed in terms of the Re_λ and the u_η . This leaves us with an expression of the form $\tau_p/\tau_T \approx \frac{1}{1+k(St/Fr)}$ as shown in Fig. 2.5(b). The predictions, shown by the solid black curve, are in reasonable agreement with our numerical observations.

While the model provides a reasonable prediction for the decrease in correlation time, the choice of u_{rms} as the characteristic velocity scale was empirical.

Simulations performed at different Reynolds numbers can answer whether this velocity scale would work for arbitrary Reynolds number. At the same time, some characteristics of the evolution of τ_p/τ_T , such as $\tau_p/\tau_T \propto (\frac{St}{Fr})^{-1}$ for large St/Fr , should be valid for all Reynolds numbers. For small and moderate St/Fr in Fig. 2.5(b), we underpredict the correlation time when compared to the simulation results. Below, we attempt to rationalize this deviation. From eq. (2.7), we note that the acceleration of a drifting particle has two contributions: (a) $D_t \mathbf{u}$ from the fluid acceleration, and (b) $\frac{St}{Fr} \partial_z \mathbf{u}$ from the velocity-gradients in the flow. In Fig. 2.6(a)–(d), we show the normalized time correlation of the particle accelerations and the gradient terms along the particle trajectories. For small St/Fr , the velocity gradient terms ($\partial_z u_z$ and $\partial_z u_x$) decorrelate slower as compared to the fluid acceleration term (see Fig. 2.6(a) & (b)). Since St/Fr is small, the fluid acceleration term $D_t \mathbf{u}$ dominates over the velocity gradient term $\frac{St}{Fr} \partial_z \mathbf{u}$. This explains the slower decrease in the decorrelation time in the numerics as compared to our predictions for small St/Fr (see $|St/Fr| < 5$ in Fig. 2.5(b)).

A few more observations may be made for the moderate St/Fr range. From Fig. 2.5(b), we note that the vertical component of particle acceleration (solid symbols) is shorter-correlated compared to the horizontal component (hollow symbols). This occurs because the fluid velocity gradient components are not all correlated in the same way. Due to the incompressibility constraint ($\partial_i u_i = 0$), the longitudinal gradient $\partial_z u_z$ is shorter-correlated compared to the transverse gradient $\partial_z u_x$, as may be seen in Fig. 2.6(a) & (b). Since the vertical acceleration is influenced by the longitudinal velocity-gradient, it decorrelates in shorter time than the horizontal acceleration. We now consider the case of large St/Fr in Fig. 2.5(b). In this case, the velocity gradient term $\frac{St}{Fr} \partial_z \mathbf{u}$ in eq. (2.7) dominates over the fluid acceleration term $D_t \mathbf{u}$, and therefore, the decrease in correlation time in DNS is in good agreement with the predictions of our eddy-crossing model.

2.4.3 Acceleration flatness

While the effect of gravity on the acceleration variance and correlation time have been comprehensively demonstrated, its role on the intermittency of particle acceleration is not clear. Intermittency, i.e. the observed strong deviations from Gaussianity, can be characterized in terms of the flatness of acceleration $\mathcal{F}(a_p) \equiv \langle a_p^4 \rangle / \langle a_p^2 \rangle^2$. Assuming statistical independence between $D_t u_i$ and $\partial_z u_i$, we obtain the tracer-normalized flatness of the particle acceleration, $\mathcal{F}(a_p)/\mathcal{F}(a_T)$, as a decreasing function of St/Fr (see Appendix A). At large St/Fr , we asymptotically approach the limits

$$\frac{\mathcal{F}(a_h)}{\mathcal{F}(a_T)} \equiv \frac{\mathcal{F}(\ddot{x})}{\mathcal{F}(D_t u_x)} \simeq \frac{\mathcal{F}(\partial_z u_x)}{\mathcal{F}(D_t u_x)}, \quad (2.15)$$

$$\frac{\mathcal{F}(a_v)}{\mathcal{F}(a_T)} \equiv \frac{\mathcal{F}(\ddot{z})}{\mathcal{F}(D_t u_x)} \simeq \frac{\mathcal{F}(\partial_z u_z)}{\mathcal{F}(D_t u_x)}, \quad (2.16)$$

It is verified that $\mathcal{F}(\partial_z u_z) < \mathcal{F}(\partial_z u_x) < \mathcal{F}(D_t u_x)$ [67]. This leads to the prediction $\mathcal{F}(a_v) < \mathcal{F}(a_h)$, i.e. vertical acceleration is less intermittent as compared to the horizontal one. In Fig. 2.5(c), we present the normalized flatness factor from our simulations for an extended $(\beta, \text{St}/\text{Fr})$ range. $\mathcal{F}(a_p)$ decreases for both buoyant and heavy particles, and the curves asymptotically reach the limits suggested by eq. (2.15) & (2.16). The accuracy of the present bubble tracking measurements is not high enough to obtain reliable estimates of the flatness. Therefore, this result is pending experimental verification.

2.5 Discussion

Our results show that the Lagrangian acceleration statistics (variance, correlation, and intermittency) of buoyant and heavy particles are very sensitive to the ratio St/Fr . To explain the origin of these in physical terms, we consider the case of a particle drifting through a turbulent flow. As the particle drifts through the flow, it meets different eddies. Owing to its short response time, the particle readjusts to the velocity of these eddies. The rate at which the particle readjusts to the new eddies is linked to the spatial velocity gradients of the turbulent flow. As a consequence, the particle experiences accelerations that the regular fluid element does not experience, thereby increasing the fluctuations (variance). The effect may be expected to become prominent when the drifting time of the particle past the energetic eddies of the flow becomes shorter than the time scale of these eddies. This explains the decrease in the decorrelation time in Fig. 2.5(b).

The same mechanism could explain the decline in the intermittency of particle acceleration. A drifting particle, instead of probing the accelerations of fluid elements, begins to sample the spatial gradients of the flow. An interesting analogy may be drawn to the intermittency of acceleration recorded by a hot-wire probe placed in a high-speed wind/water tunnel flow, where the probe effectively registers only the spatial gradients [68]. For a turbulent flow, the intermittency of the spatial gradients of velocity is lower as compared to the intermittency of the fluid element acceleration [9, 69]. Hence the observed decline in intermittency, which asymptotically approaches the expressions given by eq. (2.15) and (2.16) in the limit of large St/Fr . The above effects will be

important even for moderate Stokes number particles, and the same qualitative trends may be expected. However, a moderate St particle responds slower to the turbulent eddies it drifts through. Hence, we expect the gravity effect to be less prominent than that for the $St \ll 1$ particles we presented here.

In summary, the acceleration statistics of small Stokes number particles in turbulence is greatly modified in the presence of gravity. We report three major effects: an increase in the acceleration variance, a decrease in the correlation time, and a reduction of intermittency for buoyant and heavy particles. The ratio St/Fr governs the extent of this modification, as confirmed by our experiments using tiny air bubbles in water. Our theoretical predictions have broad validity – to particles of arbitrary density and even at large Reynolds numbers. Thus, a tiny bubble or droplet is not necessarily a good tracer of turbulent acceleration. This can be important for bubbles and droplets that drift through the turbulent oceans ($g/a_\eta \approx 100 - 1000$) and clouds ($g/a_\eta \approx 10 - 100$) [70]. On the practical side, our findings point to an important consideration when choosing bubbles or droplets for flow visualization and particle tracking in turbulent flows [43].

Wake-induced dynamics of finite-sized buoyant spheres in turbulence *

Particles suspended in turbulent flows are affected by the turbulence and at the same time act back on the flow. The resulting coupling can give rise to rich variability in their dynamics. Here we report on the acceleration statistics of finite-sized buoyant spheres in a turbulent flow. We find that even a marginal reduction in the particle's density from that of the fluid can lead to major modifications of its dynamics. The acceleration variance increases with increasing particle size, i.e in contrast to inertial range scaling predictions. We trace this reversed trend back to the growing contribution from wake-induced forces, which outweighs the spatial filtering effect for finite-sized buoyant particles in turbulence. The Galileo number (Ga) – a measure of particle buoyancy compared to viscous effects – captures the extent of this modification. At sufficiently low Ga (≤ 30) the particles respond to the carrier turbulent flow, while at higher Ga (≥ 225) their motions are dominated by wakes, with only a weak influence from the surrounding turbulent flow.

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Experiments and analysis by Mathai and Brons. Writing by Mathai, Prakash, Sun and Lohse. Supervision by Sun and Lohse, discussion of results and proofreading by everyone.

3.1 Introduction

Particulate suspensions in turbulent flows are found in a wide range of natural and industrial settings – typical examples include pollutants dispersed in the atmosphere, droplet suspensions in clouds, air bubbles and plankton distributions in the oceans, and sprays in engine combustion [9, 11, 24, 28, 58]. The behavior of a particle in a flow is intricately linked to several quantities such as the particle’s size and shape, its density relative to the carrier fluid, and the flow conditions among others [71]. For modeling purposes, the equations governing particle motion are often simplified under the point-particle (PP) approximation, which considers a dilute suspension of small rigid spheres in a non-uniform flow [61, 62]. In this framework, particle motion in turbulence is described by three fundamental control parameters: the ratio of particle size to dissipative length scale ($\Xi \equiv d_p/\eta$), the particle-fluid density ratio ($\Gamma \equiv \rho_p/\rho_f$), and the Taylor Reynolds number (Re_λ) of the carrier flow.

In many practical situations, particles have a finite size and their density can be different from that of the carrier fluid, for instance air bubbles in natural and industrial water flows [72–74], and oceanic mesoplankton that undergo vertical diel migrations [75]. In order to model such large particles in turbulent flows, one has to take into account the non-uniformity of the flow at the particle’s scale. The effects due to spatial variations in the flow across the particle have been modeled using the so-called “Faxén corrections” (FC) [24]. PP models with FC (PP-FC) have been found to work reasonably well for not too large neutrally buoyant particles [10, 13, 24, 58, 64, 76–78] and small bubbles [79] in turbulence. These studies highlighted certain effects of finite size on the particle’s statistical properties, namely a decrease in acceleration variance, an increase in correlation times, and a decrease in intermittency of the acceleration (on increasing the particle’s size). All three effects could be interpreted through classical inertial range scaling arguments, which propose spatial filtering due to increasing particle size as the underlying mechanism.

The applicability of PP-FC is questionable for finite-sized particles and in particular when the density ratio deviates from unity [80–84]. In these situations, the particle no longer adheres to one of the central assumptions of the PP models, i.e. a vanishing particle Reynolds number. For instance, when the velocity fluctuations u' at the scale of the particle are large, it can result in a finite particle Reynolds number. In another situation, the buoyancy force acting on light particles can induce finite rise velocities U_r , which can result in finite particle Reynolds numbers, $\text{Re}_p \equiv U_r d_p/\nu$. Numerical studies employing PP models have adopted a few approaches to account for Re_p effects for particles in turbulence. One involves modeling the non-Stokesian drag on the particle [64, 85]. Other methods implement two-way coupling by including

Figure 3.1: Parameter space of density ratio (Γ) *vs.* size ratio (Ξ) for particles in turbulence. Datapoints from literature: $\triangle$ - Voth *et al.* [10]; $\triangleright$ - Volk *et al.* [79]; $*$ - Gibert *et al.* [86]; $\diamond$ - Brown *et al.* [58]; $\circ$ - Qureshi *et al.* [78]; $+$ - Volk *et al.* [76]; ∇ - Mercado *et al.* [21]; $\bullet$ - Prakash *et al.* [35]; and present experiments: $\square$ - Marginally buoyant, $\square$ - Moderately buoyant and $\square$ - Very buoyant particles.

a back-reaction from the particle on the fluid [47]. But despite these efforts, there exists a lack of consensus on the fidelity of these simplified approaches.

The FC model worked reasonably well for not too large neutrally buoyant particles [24, 64, 76, 77, 79] in turbulence. This also led to its extension to predict the behavior of other classes of particles, namely heavy and buoyant particles [24, 31, 35, 79]. Building on these predictions, some generic models have been proposed to predict the *rms* of the acceleration of arbitrary-density finite-size particles [87]. These extensions practically encompass many of the naturally and industrially relevant particle-laden turbulent flows. However, in many situations, the particles do not adhere to one of the central assumptions of the PP model, i.e a vanishing particle Reynolds number. An accurate representation of particle behavior under these conditions requires fully resolved simulations (e.g. Physalis [88], Immersed boundary [77, 89, 90], Front-Tracking [91] etc), wherein one resolves the particle boundary layer in addition to a wide spectrum of turbulent eddies. However, these approaches can be computationally expensive, which restricts their feasibility to modest combinations of Re_p and Re_λ . Therefore, one is in need of experiments.

Conducting experiments with non-neutrally buoyant particles is a challenging task. In zero mean-flow turbulence setups [58, 92–94] these particles would drift vertically past the small measurement volumes, making it extremely difficult to obtain long particle trajectories in the Lagrangian frame. Hence, few experiments exist in this regime of buoyant finite-sized particles in tur-

bulence. In this chapter, we present a novel experimental strategy, wherein a mean-flow may counteract the drift of the particles. Our investigation covers the regime of finite-sized buoyant spheres ($4 \leq \Xi \leq 50$) in turbulence (see Fig. 3.1). The size-ratio (Ξ) is defined as the ratio between particle diameter d_p and the Kolmogorov length scale η in the flow. We study three density ratios ($\Gamma \approx 0.92, 0.52, \& 0.02$), where Γ is the ratio of particle density ρ_p to fluid density ρ_f . Based on the Γ value, we call the particles either marginally buoyant ($\Gamma \approx 0.92$) or moderately buoyant ($\Gamma \approx 0.52$) or very buoyant ($\Gamma \approx 0.02$). The Galileo number, $\text{Ga} \equiv \sqrt{gd_p^3(1-\Gamma)}/\nu$, provides a good estimate of the buoyancy force in comparison to viscous force. We present the Lagrangian statistics of particle acceleration, covering two orders of magnitude variation in Ga ($\approx 30 - 3000$).

3.2 Experimental setup

The experiments were performed in the Twente Water Tunnel (TWT) facility (see section 1.3), in which an active grid generated nearly homogeneous and isotropic turbulence in the measurement section [20, 21]. The water tunnel was configured to have downward flow in the measurement section, and the Taylor Reynolds number of the flow, Re_λ , was varied from 180 to 300. A small number of rigid buoyant spheres ($0.8 \text{ mm} \leq d_p \leq 10 \text{ mm}$) were dispersed in each experiment; the volume fraction ϕ was kept sufficiently low that the interaction between the spheres was negligible ($\phi \sim \mathcal{O}(10^{-5})$). The closed circuit of the water tunnel enabled the suspended particles to reappear regularly, and hence, sufficient statistics could be obtained (see Fig. 1.3).

The particles were imaged using two high speed cameras placed at a 90 degree angle between them (see Fig. 3.2). The spheres appear as dark circles in the back-lit images and their diameters corresponded to at least 10 pixels in the images. The measurement window size was adjusted to ensure this resolution, and the circle-centers were accurately detected using the Circular Hough Transform method (see Appendix B). A particle tracking code was used to obtain the trajectories of the spheres. These were further subjected to smoothing [95], which yielded robust results across the different experimental conditions (see Chapter 4 for details). Three dimensional trajectories were subsequently obtained by matching the particle tracks from the individual cameras using cross-correlation (see Movies A–D [96–99]).

Figure 3.2: A schematic of the measurement section with the two-camera 3D PTV experimental arrangement.

3.3 Results and discussion

We first address the question of how a marginal reduction in the density ratio ($\Gamma \approx 0.92$) affects the dynamics of finite-sized particles in turbulence. We focus on one of the horizontal components. In Fig. 3.3(a), we show the normalized acceleration probability density function (PDF) for marginally buoyant (MB) spheres at different Ξ . The Re_λ is maintained constant (≈ 180), and the particle size ranges from a few Kolmogorov lengths to a fraction of the integral scale. At the smallest size ratio ($\Xi \approx 4$) the PDF has wide tails, while it is nearly Gaussian at $\Xi \approx 50$. This trend of narrowing PDF tails with increasing Ξ was predicted for buoyant particles using the FC model [24, 35]. Similar behavior was reported experimentally and also predicted numerically (FC model) for neutrally buoyant spheres in turbulence [24, 76].

While it seems that accounting for flow non-uniformity and spatial filtering are sufficient to capture the narrowing of the tails of the PDF, the absolute accelerations reveal something very different. The inset to Fig. 3.3(a) shows that the *rms* of particle acceleration ($\langle a^2 \rangle^{1/2}$) increases with Ξ , suggesting

Figure 3.3: Acceleration statistics for buoyant particles in turbulence at $\text{Re}_\lambda \approx 180$. (a) Acceleration PDF for marginally buoyant spheres. Inset shows $\langle a^2 \rangle^{1/2}$ (mm/s²) vs. $\Xi \equiv d_p / \eta$. The dash-dotted line (in blue) with slope $\approx -1/3$ shows the prediction based on classical inertial scaling arguments [78]. (b) Normalized acceleration variance from experiments (EXP) compared to the results from Faxén-corrected (FC) simulations at $\text{Re}_\lambda = 180$ [24]. The horizontal dashed and dash-dotted lines (in black) are lines marking the tracer particle acceleration ($\langle a^2 \rangle / \langle a_f^2 \rangle = 1$) and the upper bound of FC simulations ($\langle a^2 \rangle / \langle a_f^2 \rangle = 9$), respectively. Here, a_f is the fluid tracer acceleration.

that larger particles experience more intense accelerations. This observation is in contrast with both inertial range scaling predictions ($\langle a^2 \rangle^{1/2} \propto \Xi^{-1/3}$) and the predictions of FC model (Fig. 3.3(b)), according to which larger particles should experience milder acceleration fluctuations due to spatial filtering effects [10]. While the $-1/3$ scaling has been experimentally validated for neutrally buoyant (NB) spheres in turbulence [10, 58, 78], interestingly here, only a marginal reduction in particle density reverses the trend. Furthermore, the accelerations exceed even the upper bound of the Faxén model predictions for buoyant particles i.e. $\langle a^2 \rangle / \langle a_f^2 \rangle > 9$ (see marginally buoyant ($\Gamma \approx 0.92$) case in Fig. 3.3(b)). Here, a_f is the measured fluid tracer acceleration, which was calculated from particle tracking experiments.

Figure 3.4: (a) Lagrangian acceleration autocorrelation function $C_a(\tau)$ for marginally buoyant (MB, $\Gamma \approx 0.92$) and neutrally buoyant (NB, $\Gamma \approx 1$) spheres with low Ga (≤ 30), along with fully resolved DNS predictions from Ref. [77] for NB particles at $Re_\lambda = 32$. (b) $C_a(\tau)$ for MB particles ($\Gamma \approx 0.92$) with moderately large Ga ($225 \leq Ga \leq 900$) at $Re_\lambda \approx 180$.

In order to explain the deviations, we look into the temporal response for the marginally buoyant particles. The Lagrangian autocorrelation function of acceleration $C_a(\tau)$ plotted for different particle size-ratios are shown in Fig. 3.4(a) & (b). At lower size-ratios (see Fig. 3.4(a)), the particle accelerations decorrelate according to a turbulent spatial filtering based time scale, $\tau_d = (d_p^2/\epsilon)^{1/3}$. This corresponds to the MB particle with the smallest size-ratio ($\Xi \approx 4$) in the inset to Fig. 3.3(a). At larger size-ratios (see Fig. 3.4(b), where $\Xi \approx 16, 25$, and 50), the accelerations decorrelate according to a vortex shedding time-scale, $\tau_v = \frac{d_p}{St \times U_r}$, instead of τ_d . Here St is the Strouhal number, and U_r is the measured mean drift velocity of the particle in the turbulent flow. In this regime, the decorrelation timescale is nearly constant, with only a weak dependence on the turbulence (or Re_λ .) We therefore see a transition from the steady drag regime ($Ga \approx 30$) to a regime with growing

vortex-shedding-induced effects (Ga from 225 to 900).

The autocorrelation functions revealed the dominant factors governing the dynamics of the buoyant particles in our experiments. However, this cannot explain the observed increase in the acceleration variance. For this, we review the two contributions which can lead to a change in $\langle a^2 \rangle^{1/2}$: (i) the contribution from turbulent fluctuations - this decreases with increase in particle size due to spatial filtering, and (ii) the contribution from vortex-shedding - this increases with increasing particle size. Therefore, the total $\langle a^2 \rangle^{1/2}$ is the combined effect of turbulence-induced and vortex-shedding-induced forcings. As Ξ increases from 4 to 50, the unsteady forcing due to vortex shedding starts to outweigh the turbulent forcing at the particle's scale. This results in an overall increase in $\langle a^2 \rangle^{1/2}$ with Ξ . It is interesting to note that despite the increase in the $\langle a^2 \rangle$ with Ξ , the normalized acceleration PDFs (in Fig. 3.3(a)) showed a trend of narrowing tails with increasing Ξ . However, it is important to note that this trend of decreasing intermittency (i.e narrowing of the tails of acceleration PDFs) seen here and in prior investigations [35] is rather a coincidence, since the absolute particle accelerations are orders of magnitude larger.

Next we study the effect of the density ratio Γ on buoyant sphere acceleration statistics. In Fig. 3.3(b), we show the results of varying particle density at constant $\Xi \approx 50$. Not surprisingly, the FC model underpredicts the acceleration variance for all three density ratios. The deviation is greatest for the ‘very buoyant’ particle ($\Gamma \approx 0.02$; Ga ≈ 3000), whose acceleration is almost three orders of magnitude higher. In Fig. 3.5(a) & (b), we present the horizontal and vertical components of the velocity and acceleration PDFs for this very buoyant particle at three Re_λ . The horizontal velocity PDF shows a symmetric *flat-head* distribution, and the horizontal acceleration PDF has a *bumped-head* distribution. Both the velocity and acceleration statistics show no observable effect of changing Re_λ or Ξ . The vertical components have similar characteristics, but the distributions show positive skewness, particularly for acceleration. The Lagrangian autocorrelations of velocity and acceleration are oscillatory (see Appendix B), with a period comparable to τ_v .

Figure 3.5: PDFs of (a) horizontal and (b) vertical components of velocity and acceleration for a *very buoyant* ($\Gamma \approx 0.02$) finite-sized sphere at $Ga \approx 3000$. Blue, green, and red colours correspond to $Re_\lambda \approx 180$, 250, and 300, respectively. The variable X corresponds to v_h , a_h , v_v and a_v in their respective PDFs. The PDFs have been shifted vertically for clarity in viewing. A Gaussian profile (thin black line) is overlaid for comparison. (c)-(f) show time traces of velocity and acceleration components of a representative trajectory at $Re_\lambda \approx 300$. (g)-(j) show the trajectory corresponding to the time traces in (c)-(f), projected on a vertical plane and color-coded with instantaneous quantities.

In Fig. 3.5(c)-(f), we plot the time-series of velocity and acceleration of a representative particle trajectory of the ‘very buoyant’ sphere from our experiments. A careful examination reveals that the non-regular PDF shapes have their origin in the time history of the cyclic motions undergone by the particle. In particular, the *flat-headed* velocity, the *bumped-head* acceleration, and the positively skewed vertical PDFs can all be traced back to the specifics of the periodic motions. In Fig. 3.5(g)-(j), we plot the trajectory corresponding to this time trace projected on a vertical plane. The particle traces regular *Lissajous* orbits, only weakly disturbed by the surrounding turbulence (see Movie D [99]). While such periodic motions are classically observed for tethered spheres in laminar flows [25, 100], it is interesting to observe these in a turbulent flow.

In summary, the dynamics of finite-sized buoyant particles in turbulence can be strongly two-way coupled. We track the movements of spheres with a marginal density difference ($\Gamma = 0.92$; $4 \leq \Xi \leq 50$) in turbulence ($\text{Re}_\lambda = 180 - 300$). We show that, despite including the finite-size corrections, the point particle models break down in situations where the particle density is not perfectly matched with the fluid. The effects of finite particle Reynolds number (Re_p) are evident in the experiments. In the present experiments, this occurs due to the finite rise speed of the particles, which causes the forcing due to vortex shedding to outweigh the forcing due to turbulence at the particle’s scale. The same could be true in intensely turbulent environments where the turbulent forcing outweighs particle wake effects. In these flows, the large velocity fluctuations (u') at the particle’s scale could lead to finite $\text{Re}_p \equiv u'd_p/\nu$ [29].

Tracking the translation and rotation of a large buoyant sphere in turbulence

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We report experimental measurements of the translational and rotational dynamics of a large buoyant sphere in a turbulent flow. We present a method to simultaneously determine the position and (absolute) orientation of the sphere from visual observation. The method employs a minimization algorithm to obtain the orientation from the 2D projection of a specific pattern drawn onto the surface of the sphere. This has the advantages that it does not require a database of reference images, is easily scalable using parallel processing, and enables accurate absolute orientation reference. The translational accelerations are anisotropic and show evidence for periodic motions. The angular autocorrelations show weak periodicity. The angular acceleration PDFs exhibit wide tails, however without a direction dependence.

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4.1 Introduction

Understanding particle-laden flows is crucial to make predictions on global phenomena of interest such as pollutant transport and mixing in industrial processes ([9]). The particles carried by such flows can be large in size, and their density is usually different from that of the surrounding fluid. Theoretical studies often model such objects as passive, finite-sized particles advected by the flow ([24, 61, 101]). However, as we have seen in Chapter 3, these models are limited in their applicability to vanishing particle and shear Reynolds numbers. Often, the density-mismatch of particle with the fluid causes finite rise velocities, which leads to finite particle Reynolds numbers ([102], [103], [104] etc). Much variability can be expected in the dynamics of such particles in flows ([22, 100, 105, 106]). The forcing responsible for these varied dynamics is linked to the vorticity shed in the wake of the particle ([22, 107, 108]). However, these forces may not act along the geometric centre of the particle. Therefore, it is possible that these could as well induce torques on the particle. Little is known about the resulting translational and rotational dynamics, particularly for buoyant particles in turbulent flows. This motivates us to develop a reliable measurement technique for tracking the translation and rotation of a spherical particle in a flow.

In three dimensions, an object's location and orientation can be fully described by six independent variables. Many physical experiments rely on image analysis to obtain these parameters from experimental data ([109, 110]). Most of these systems capture translation and retrieve orientation from the relative motion of translating nodes ([81]). In the more elementary forms, the translation of nodes could be used to determine the velocities of the particles. However, these methods were not accurate enough when higher derivatives of orientation were to be determined. Recently, Zimmermann *et al.* [111] introduced a method based on the identification of possible orientation candidates at each time step using projections of a pattern painted on the surface of a sphere. They found surprisingly intermittent behavior in the acceleration statistics of a neutrally buoyant sphere of diameter of the order of the integral scale.

In the present work, we introduce buoyancy to the problem of spherical particle dynamics in turbulence. We recover the translation and rotation as a function of time. The core of the method, which is to compare experimental images to synthetic ones is the same as proposed by [111]. The novelty here lies in the way the pattern is generated (both physically on the surface of the particle and numerically for the comparison with actual images to match the orientation). Hence, the synthetic images for any given orientation are analytically known and do not need to be determined from static images.

Furthermore, the method is easily scalable using parallel processing. These aspects are elucidated in section 4.3. In addition, we describe a smoothing spline based methods that enables accurate representation of the higher derivatives of experimental data. Finally, in section 4.5, we present the experimental results.

4.2 Experimental Setup

The experiments were conducted in the Twente Water Tunnel facility (TWT), designed to study particle-laden flows (see Fig. 1.3 and [112]). The measurement section has dimensions $0.45 \times 0.45 \times 2 \text{ m}^3$, with three glass walls providing optical access to perform particle tracking measurements. The setup houses an active grid above the measurement section, consisting of 24 independently rotating motors, which produce nearly homogeneous and isotropic turbulence with Re_λ up to 300 in the downstream section of the water tunnel. The flow in the measurement section was characterized using a cylindrical hot film probe (Dantec 55R11) by following the same methodology as reported in [21]. The experiment reported here was performed at $Re_\lambda \approx 300$. The dissipation rate $\epsilon = 505 \times 10^{-6} \text{ m}^2/\text{s}^3$, and the dissipation length and times scales were approximately $211 \text{ }\mu\text{m}$ and 44 ms respectively. The sphere used in this study has a diameter, $d_p = 25 \text{ mm}$, with an effective mass ratio, $m^* \approx 0.82$, where m^* is the ratio of the mass of sphere to the mass of the sphere's volume of water [108]. The mass ratio was chosen such that the mean rise velocity of the sphere matched the mean downward flow velocity in the measurement section. This was necessary for obtaining sufficiently long sphere trajectories and for converged statistics. The spheres were designed as spherical shells with the MP300 resin, with a bulk density $\rho_b \approx 1089 \text{ kg/m}^3$. The spherical shells were made using a 3D printing technique, and the surface roughness was within 50 microns. The printing resolution depended on two factors. First, the resolution of the 3D printer used for making the stencil. In the present case, we could produce complex stencil designs with patterns as small as a millimetre on a 25 mm sphere. The second limiting factor was the painting procedure itself. Once the stencil was fixed on the sphere, the pattern had to be spray painted. Here again, it was practical to spray paint patterns of 2 mm dimension.

The recordings were made with two Photron PCI-1024 high-speed cameras at 500 fps and megapixel resolution (see Fig. 4.1(a) & (b)). The cameras were positioned at a 90° angle between them and focused at the center of the test section on a $150 \times 150 \text{ mm}^2$ area, which resulted in a spatial resolution of $150 \text{ }\mu\text{m}/\text{pixel}$. The images showed that perspective effects were negligible.

Figure 4.1: (a) Measurement section of the Twente water tunnel with orthogonal camera experimental arrangement. A painted sphere (not to scale) is shown in the figure, which is viewed by both cameras. (b) The sphere and its trajectory as recorded by one of the cameras.

The measurement volume was lit by eight 20 W LED lamps from the sides. The flow velocity was tuned to ensure that the spheres stayed in the viewing window for a considerable duration.

4.3 Method

The position and orientation of the sphere are determined using image analysis. The background was painted gray for contrast between the black and

white pattern on the sphere. The sphere is separated from the background by subtracting the absolute difference in intensity from the background image. Subtracting the absolute difference yields a full dark circle with a brighter background. The sphere is then detected using the Circular Hough transform (CHT) technique and the centroid of the detected circle gives its position. The accuracy of CHT varied with angle of view. Therefore, choosing a good pattern is one of the most important steps. We found that a pattern that contained almost equal fraction of dark and bright areas for most angles of view improved the circle detection accuracy. Additionally, we used the outputs from previous frames to improve and speed up the detection process. The CHT also returns the detected circle diameter D_c in pixels from the image. Since the sphere diameter (25 mm) is known, we use a calibration $(25/D_c)$ mm/pixels for the recorded images. This ensures a correction for the minor magnification changes due to the sphere moving forward or back in the measurement section.

There are several approaches to obtaining the three dimensional trajectory of a particle. One approach is to reconstruct the spatial positions using multiple cameras ([113]). Precise spatial reconstruction is needed when there are multiple particles in the image and identifying and separating particles can be a challenge. In the present work, there are only a few particles in the measurement volume. We use a different method using two orthogonal cameras to obtain the three dimensional trajectory. The camera arrangement is such that the magnification and field of view are comparable between the two cameras. The camera field of view covered the full width of the water tunnel. We first obtain two dimensional trajectories of particles from both cameras. The redundant data corresponding to the vertical motion for two cameras is used to compare individual trajectories. In order to avoid ambiguities, we cross-correlate the vertical acceleration time series from the two cameras. A match is found when the cross-correlation coefficient between the two acceleration time series is greater than 0.98. This was done even for long trajectories, such as the one shown in Fig. 4.1(b). This yields a simple yet robust three dimensional track of the particles without needing to resort to complex spatial reconstruction algorithms.

Obtaining the rotation of the sphere is a more complex task compared to position tracking. Hence, the remainder of this section will be used to describe the method that is used to obtain the orientation. The method to determine the orientation can be divided into four parts. Initially, a suitable boolean surface pattern is created. This pattern can be described as a piecewise constant analytic function $F(\theta, \phi)$ such that, given a coordinate on the surface of the sphere, the function returns either a zero or a one depending on the color of its corresponding infinitesimally small surface element. Here, θ and ϕ are

Figure 4.2: (a) 3D stencil used for painting the pattern on the sphere. (b) Projection of the pattern on the sphere onto a plane.

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the azimuthal and the polar angles respectively, and $F(\theta, \phi)$ is radius independent. Second, the pattern is drawn onto a physical sphere. This is realized using a 3D-printed painting stencil and an airbrush system (see Fig. 4.2(a)). It is imperative that $F(\theta, \phi)$ is painted as accurate as possible to decrease the introduced error in this step. Third, a synthetic 2D image is constructed from a projection of the surface of the sphere onto a plane (see Fig. 4.2(b)). The synthetic image for any given orientation is analytically known and does not need to be determined from static images. The projection is a function of the angle of rotation of the sphere and can be conceptually understood as the analog of the projection of the physical sphere on the recording camera. Finally, the rotated and projected synthetic pattern is compared to an image of the physical sphere. The minimizing function can take different forms. For instance, a cross-correlation function between the synthetic pattern and the camera image could be used to find the best match. Alternately, a suitable cost function may be used to search for a match. Here we use a cost function, defined as the sum of the absolute difference between the binarized image pixels and the corresponding pixels in the synthetic image. The orientation for which this comparison yields the best match is then determined using a Nelder-Mead minimization algorithm. These steps are illustrated in a flowchart in Fig. 4.3. Four reference orientations are shown in Fig. 4.4(a). A sample set of experimental images and synthetically recovered orientation output by the method are given in Fig. 4.4(b) & (c). The Fortran90 code and the stl-file of the stencil used for painting the spheres can be found in the supplemental material of [59].

The choice of painted pattern is an important step in the method of ori-

Figure 4.3: Orientation detection method flowchart

Figure 4.4: (a) Several viewpoints of the synthetic pattern. Here, the synthetic projection is defined in the axis-angle convention as front: $\rho = (1, 0, 0, 0^\circ)$, left: $\rho = (0, 0, 1, -90^\circ)$, top: $\rho = (0, 1, 0, -90^\circ)$ and back: $\rho = (0, 0, 1, 180^\circ)$. (b) Experimental images of the sphere with the painted pattern for arbitrary orientations. (c) Synthetic image equivalents found by the orientation detection program corresponding to the experimental images in (b).

Figure 4.5: PDFs of the difference between actual orientation and orientation as found by the orientation detection method for different Signal-to-noise ratios (SNR). a) $\text{SNR} = \infty$ and b) $\text{SNR} = 2$. Here N is the width of the sphere in the image in pixels. N takes on values 50, 70 and 100 in the three cases shown. With increasing resolution N , the width of the PDF decreases.

entation detection. A necessary condition is that the projection of $F(\theta, \phi)$ onto a plane must be unique for each orientation of the sphere. Additionally, it is desirable that the pattern contains a minimum of edges and corners. The latter criterion is essential for fast convergence of the algorithm to the global minimum. Here we use the axis-angle method to describe the orientation considering its straightforward and singularity-free definition. It allows for smooth and continuous rotation from any orientation including around the Euler-angle singularities as it does not suffer from gimbal lock problems.

The pattern $F(\theta, \phi)$ represents a simply connected region, even then, some local minima may arise due to the two-tone color limit. When the orientation is unknown, for instance in the first frame of a movie, a comparison of several

Figure 4.6: Numerical error scaling for signal-to-noise ratio $SNR = \infty$ and $SNR = 2$.

initial orientation estimates solves this local minima problem. This yields the global minima. However, every initial estimate reduces the performance. Fortunately, given a sufficiently high ratio of frame rate over rotation rate, any frame in a sequence may use the outcome of the previous frame as initial estimate, reducing the number of initial estimates required to just one. This *causal* method is the preferred method for analyzing large sequential datasets. The computational performance of the method is approximately 30 frames per second on a contemporary computer, suggesting that processing of large datasets is straightforward.

To determine the numerical accuracy of the method, several sets of 1024 projections are created using pseudo-random generated orientations. These synthetic images are used as input to the algorithm. A probability density function of the difference between the actual orientation and the orientation as found by the algorithm is shown in Fig. 4.5(a) & (b), and the corresponding scaling of the standard deviation as function of the image width N (in pixels) is shown in Fig. 4.6. The standard-deviation of error scales as N^{-2} . Assuming a Gaussian error distribution, this means that less than 1% of the measured data showed an error larger than 1 degree for $N \geq 50$. Also, Fig. 4.6 shows that noise decreases the accuracy to some extent but does not affect the reliability of the algorithm. However, in the actual experiment, most image artifacts arise due to shadows or glare. These may be reduced by using diffuse light-sources and a matte-paint finish on the surface of the sphere.

The axis-angle output, $\rho \equiv (k_x, k_y, k_z, \alpha)$, is defined with respect to a reference orientation in the camera coordinate system (see leftmost image in Fig. 4.4 (a)). A typical output obtained for the case of a buoyant sphere in a turbulent flow is given in Fig. 4.7(a). This output has little physical signifi-

Figure 4.7: (a) Axis-angle output obtained for a long trajectory of a buoyant sphere in a turbulent flow. (b) Orientation $(\theta_x, \theta_y, \theta_z)$ obtained from Camera 1. (c) Absolute difference between independent orientation measurements from Camera 1 and Camera 2. The maximum deviation is within 2.5 degrees.

cance. Quantities of relevance to a rotating sphere are its angular velocity ω and angular acceleration α , which are directly related to the rotational kinetic energy and the torque acting on the sphere. We adopt the following approach to compute ω and α in the lab coordinate system. Firstly, a high frame rate is used for accurate estimation of the time derivatives. In the present case, a frame rate of 500 fps ensures about 30 recordings in one Kolmogorov time. Therefore, within the inter-frame time interval $\Delta t = 1/500$ sec, the particle's angular velocity may be assumed constant.

According to Euler's theorem, the angular velocity is linked to the inter-frame axis-angle, $\rho_\Delta \equiv (k_{\Delta x}, k_{\Delta y}, k_{\Delta z}, \Delta\alpha)$, by the relation

$$\vec{\omega}(t) = \frac{\Delta\alpha}{\Delta t} (k_{\Delta x} \hat{i} + k_{\Delta y} \hat{j} + k_{\Delta z} \hat{k}) \quad (4.1)$$

Now ρ_Δ may be obtained from the reference orientation based axis-angle output (ρ) through a few transformations. Firstly, the axis angle output is converted to rotation matrix using Rodrigues formula. Let $[R_{0,i}]$ and $[R_{0,(i+1)}]$ represent the rotation matrices corresponding to rotations from the reference orientation to the orientations in i^{th} and $(i+1)^{th}$ images respectively. Then $[R_{0,i}] \times [R_{i,(i+1)}] = [R_{0,(i+1)}]$ and consequently, the inter-frame rotation matrix $[R_{i,(i+1)}] = [R_{0,i}]^T \times [R_{0,(i+1)}]$. This follows from the identity that $[R]^{-1} = [R]^T$ for rotation matrices. From $[R_{i,(i+1)}]$, $\rho_\Delta \equiv \rho_{i,(i+1)}$ may be obtained, and eq. 4.1 gives the angular velocity $\vec{\omega}(t)$ in the lab coordinate system.

The accuracy of the detection in a real experiment needs to be verified separately. For this, we employ a two-camera arrangement. If the orientations determined by the two cameras are comparable, then the method can be regarded accurate. Fig. 4.7(b) shows the three orientation angles obtained from

Camera 1 for a long particle trajectory. The motion of the same sphere was captured by Camera 2, which was placed at a 90° angle with Camera 1. The difference in orientation prediction is plotted in Fig. 4.7(c). The maximum deviation in this recording is within 2.5 degrees. These differences are due to experimental noise and may be filtered out in the data smoothing step, to be explained in the following section.

4.4 Evaluating higher derivatives from experimental data

Evaluating the forces and torques on a body moving through a flow is one of the key steps in understanding the dynamical system. One way to do this is by using accelerometers. However, this may not always be convenient since a direct measurement of acceleration usually requires attaching devices to the body. Moreover, the method is intrusive in nature and often leads to variations in the mass and center of gravity of the system ([114]). In such situations, it may be suitable to numerically compute the acceleration from the second derivative of position of the body. This approach is the basis of our present study, where we determine both the position and orientation of a large buoyant sphere in turbulence using recordings from a high speed camera.

Estimating the derivatives from experimentally determined position and orientation data is a nontrivial task. The difficulty arises because the data obtained from experiments is bound to have inherent noise due to measurement uncertainties. Historically, there have been two popular methods for smoothing particle trajectories in turbulent flows. One method involves fitting parts of the particle trajectory to a polynomial of second order or higher. Voth *et al.* [10] used a second-order polynomial, Lüthi *et al.* [115] and Martínez-Mercado *et al.* [21] used a third-order polynomial. Other researchers have used a Gaussian kernel for smoothing ([76, 116]). Both the methods employ piecewise discontinuous fitting of analytical functions to smooth out the noise. The effectiveness of these methods in filtering out the experimental noise depends on the fitting parameters chosen. In spite of the extensive literature on smoothing methods, there still prevails a general lack of consensus on the method of finding the optimal fitting parameter. Existing guidelines for the choice of the fitting parameters, such as those suggested by [116], require knowledge of the smallest time scales in the flow *a priori*, and therefore are likely to introduce a bias into the analysis.

In this chapter, we explore an alternate method using smoothing splines to filter out the experimental noise. The method is based roughly on the work by Epps [117]. While this was originally used to estimate the translational accelerations experienced by a water-entering object ([95, 118]), the method

Figure 4.8: Roughness vs error frontier for a family of candidate smoothing splines. The kink marks the critical error tolerance of the optimal smoothing spline. Inset shows $\log(C(s))$ vs. $\log(E(s))$. The optimal fit corresponds to the maximum curvature point indicated by the red star symbol.

4

may be applied to particle trajectories in turbulence. Consider a general set of experimental data $y(t)$ acquired at high temporal resolution. We use a smoothing spline based roughness limiting method to reduce the experimental noise in the data. The function *spaps* in Matlab requires mainly the error tolerance as the input parameter: $E(s) = \int_{t_1}^{t_N} |y(t) - s(t)|^2 dt$. We use a roughness estimate, $R(s) = \int_{t_1}^{t_N} |\frac{d^3 s}{dt^3}|^2 dt$, to scan for the most suitable fit to the experimental data. This is based on the assumption that the true function does not have very large changes in acceleration. Large fluctuations in the acceleration are expected to arise from noise in the data. In Fig. 4.8, we demonstrate the step to determine the optimal fit from experimental data. Increasing the error tolerance beyond E_{crit} does not reduce the roughness of the curve. For error tolerances below E_{crit} , the roughness of the fits are increased significantly, indicating that the noise is retained. Therefore, the fit corresponding to the kink can be thought to best represent the true curve.

We test the sensitivity of the velocity and acceleration to the choice of the spline. For this, we use a sample analytical function defined as a combination of sine functions, $y = a \sin x + b \sin 2x + c \sin 4x + d \sin 10x$, where a , b , c and d were generated randomly in the $[0; 1]$ range, and synthetically introduce random noise in the signal (see Fig. 4.9(a) and inset). In Fig. 4.9(c)–(h), we compare the velocity and acceleration estimates for critical (E_{crit}), sub-critical (E_1) and super-critical (E_2) error tolerances. The sub-critical error tolerance (Fig. 4.9(f)) yields very high accelerations, while the super-critical error tolerance over-smoothes the curve (Fig. 4.9(h)). The acceleration esti-

Figure 4.9: (a) Analytical function and the noisy data with $\text{SNR} \approx 4$. Inset shows a zoom-in of the noisy data. (b) Roughness vs error plot for the noisy data. First and second derivatives (velocity and acceleration) estimated for (c) & (d) $E = E_{crit}$, (e) & (f) $E = E_1$, and (g) & (h) $E = E_2$. $E = E_{crit}$ gives an accurate estimate of even the second derivative, except at the endpoints.

mate for critical error tolerance compares best with the second derivative of the true function.

Another interesting feature of the present method is that the kink (fig-

ures 4.8 & 4.9(b)) flattens out for experimental data with low signal-noise-ratios. Thus, the method also serves as a check for the quality of data, which can be difficult to estimate for standard smoothing methods. Alternately, the spline fit guess for the particle acceleration from a few representative tracks could be used to determine the optimal parameters for Gaussian Kernel smoothing ([119]). This can speed-up the fitting process when sampling large datasets, while ensuring that the fitting windows are properly chosen. This was the approach adopted in Chapter 2.

4.5 Experiments and discussion

We present some experimental results on the translational and rotational motion of a large buoyant sphere in nearly homogeneous and isotropic turbulence. We first address the question of how the particle's translational velocity and acceleration decorrelate in time. In Fig. 4.10 (a), we plot the Lagrangian autocorrelation function for the horizontal (x & z directions) velocity and accelerations. The particle response scale is fairly well predicted by the relation $\tau_v = d_p / (St \times U_r)$, where St —Strouhal number ≈ 0.2 , and U_r is the measured mean rise velocity relative to the flow. Both the velocity and acceleration decorrelate in the same time ~ 0.5 sec, and they both display periodicity. This behavior is similar to the results in Chapter 3.

Fig. 4.10(b) shows the angular velocities and accelerations in the horizontal direction. The time axis is normalised with the particle-sized eddy time scale, $\tau_d = (d_p^2 / \epsilon)^{1/3} \sim 1$ s. The angular acceleration crosses the first minima in approximately $1 \tau_d$, which is four times longer compared to the minima crossing time for translational acceleration. The angular velocity components decorrelate even slower than the angular acceleration, and they show only a weak periodicity. Surprisingly, we do not find any preferential orientation for the angular velocity vector with the translational acceleration vector, unlike the findings of Zimmermann *et al.* [120]. Therefore, in the present case, particle rotation does not appear synchronized with its translational motion, and we find no evidence suggesting a lift force. These aspects will be discussed in more detail in Chapter 5.

In Fig. 4.11(a) & (b), we show the probability density functions (PDFs) of the horizontal and vertical components of the particle velocity and acceleration. The horizontal velocity PDF shows a symmetric flat-head distribution with sub-Gaussian tails. The behaviour is notably different from the nearly Gaussian velocity PDFs found for neutrally buoyant particles ([21]). The flat head of the PDF may be explained from a typical time series of velocity of the sphere. This is similar to the observations made for a very buoyant

Figure 4.10: (a) Lagrangian autocorrelation function of (a) translational and (b) rotational velocities and accelerations for the buoyant sphere in turbulence.

sphere ($m^* \equiv \rho_p/\rho_f \approx 0.02$) in Chapter 3. Here again, the particle undergoes repeated cyclic motions, weakly disturbed by the carrier flow. The velocities in these cycles lie in the $\pm 1.5 \langle a^2 \rangle^{1/2}$ range. More extreme accelerations are less frequent and contribute to the low probability tails of the PDFs. The horizontal acceleration PDF is Gaussian, which again is similar to the observations for the $\Xi \approx 50$ sphere in Chapter 3. The vertical velocity and accelerations are positively skewed distributions, clearly indicating a directional dependence. This strong anisotropy is not inherent in the carrier turbulent flow [21], and therefore, must originate from the up-down asymmetry set up by the vortex-shedding behind the sphere. The PDFs of the rotational quantities reveal a different story (Fig. 4.11(c) & (d)). Both the horizontal and vertical angular velocities follow a nearly Gaussian distribution. The angular acceleration PDFs show symmetric distributions, but with wider tails than the translational acceleration PDFs. It may be noted that Re_λ and particle size ratio $\Xi \equiv d_p/\eta$ are comparable to the experiment of [111]. Thus, changing only the particle-fluid density ratio brings in these new effects, leading to observable

Figure 4.11: Velocity and acceleration PDFs in the (a) transverse and (b) streamwise directions for the buoyant sphere (25 mm diameter). Angular velocity and angular acceleration PDFs in the (c) transverse and (d) streamwise directions for the sphere. The PDFs have been shifted vertically for clarity of viewing. The black solid and dotted curves are Gaussian distributions, given for comparison.

differences in the particle's dynamics.

4.6 Conclusions

We have conducted experimental measurements of the dynamics of a large buoyant sphere in a turbulent flow and described the methods and procedures to track simultaneously its position and orientation. A fast, accurate and adaptable method that determines the absolute orientation of the sphere in 3D space is described and validated. The standard deviation of error for this method is approximately 0.2 degrees. Distortion using salt & pepper noise with SNR= 2 increases the standard deviation of error to 0.7 degrees, however without impacting the reliability of the method. We followed the approach of using the experimentally determined data to calculate time derivatives (velocity and acceleration) of position and orientation. To this end, we employed a smoothing spline based method, which enables accurate estimation of higher derivatives. The method has the advantage that it yields an optimal fit that is continuous and differentiable. The six-degree of freedom particle tracking, in

combination with the smoothing method introduced here offers a robust way of obtaining a complete description of the dynamics of a spherical particle in a turbulent flow. The method will be extensively used in Chapter 5.

Controlling particle path-instabilities by tuning rotational inertia *

Particles suspended in fluid flows display various kinds of path-instabilities, which depend on their interactions with the flow. The conditions which onset these oscillatory motions have remained difficult to predict even for simple bodies, such as spheres and cylinders rising through a fluid. In literature, two parameters have been considered to control the buoyancy-driven motion of a spherical particle in a flow; they are the particle's density relative to the fluid ($\Gamma \equiv \rho_p/\rho_f$) and its Galileo number (Ga). Consequently, once these are specified, the opportunities for variation in particle dynamics appear limited. Here we show that path-instabilities can be triggered by reducing the particle's moment of inertia. We reveal this in two canonical flow environments: a) a turbulent flow, and b) in quiescent liquid. In the turbulent flow, this leads to an enhancement in the particle's translational accelerations and a transition from a tumbling to a fluttering state of motion. In quiescent liquid, reducing the moment of inertia induces a modification of the particle wake, resulting in large-amplitude path-oscillations. Thus, rotation can couple with a particle's translational motion, triggering modifications to its wake-induced oscillatory dynamics. This new perspective might also explain some of the variability that

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Experiments, analysis, and writing by Mathai and Zhu. Supervision by Sun and Lohse, discussion of results and proofreading by everyone.

has been witnessed in the dynamics of buoyant spheres and cylinders.

5.1 Introduction

The buoyancy-driven motion of particles rising or falling through a fluid has captured the interest of scientists and engineers for centuries. From the earliest observations of rising bubbles by Da Vinci to the falling hog-bladder tests by Newton, the variability observed in the trajectories of these particles has been intriguing [105, 121]. Such variability can have far-reaching implications for many natural and industrial particle-laden and tethered-body flows, wherein the movement of the particles can significantly alter the drag, and the transport of heat and nutrients in the fluid [122–126]. For example, the presence of rising bubbles and particles in the oceans results in warmer ocean waters and larger positive temperature gradients [7, 8]. Similarly, in industrial reaction catalysis, buoyant particles or bubbles are often released to enhance the mixing in the fluid [127, 128]. From these observations, a link can be inferred between the oscillatory motions and wakes of the rising particles, and the processes of turbulent diffusion at work [129, 130]. Several studies have emphasized the importance of this [131–134]; however, few have established the factors which trigger the oscillatory motions for buoyant particles in fluids.

The motion of a buoyant particle in a flow is a complex two-way coupled problem. The particle moves through the fluid in response to the fluid forces, and this motion in turn exerts a back-reaction on the flow. The path oscillations that result, have been found to be robust to interactions between neighbouring particles, and to changes in the flow conditions, making studies of isolated bodies relevant to various multiphase flows [135, 136]. In order to model such systems, often a prototype problem is considered. The particle is modeled as a sphere, and the flow is assumed to be turbulent, homogeneous, and isotropic. For the particle dynamics, two control parameters have been proposed to be of relevance: the particle to fluid density ratio Γ , and the particle Galileo number Ga expressing the ratio between gravity, buoyancy and viscous effects [105]. Among these, Ga governs the onset of various kinds of wake-instabilities behind the particle, and Γ governs the movement of the particle in response to these wake-instabilities and vortex-induced forces. Once the background turbulence is also considered, two additional parameters need to be included: the particle's size in relation to the dissipative length scale of the flow ($\Xi \equiv d/\eta$), and the Taylor Reynolds number Re_λ of the flow [9]. Hence, knowledge of these four dimensionless groups: $[\Gamma, \Xi, Re_\lambda, Ga]$, along with Fr (see Chapter 2) would seem to define the problem. Expecting these dependences, researchers have performed extensive studies on the dynamics of light and heavy particles in a range of flow environments [57, 58, 78, 100, 137–140].

A detail that has often been overlooked is the rotational dynamics of buoy-

Figure 5.1: (A) A perspective representation of the measurement section of the Twente Water Tunnel (TWT) with the two camera experimental arrangement for tracking the trajectory and orientation of buoyant spheres in turbulent flow (not to scale). (B) & (C) Cut-section schematics of the high and the low moment of inertia (MoI) spheres, respectively. The high MoI sphere is hollow. The low MoI sphere is constructed with a dense metallic core, which is held in the center using thin spokes as shown in (C).

5

ant spherical particles. Theoretical studies have mostly restricted their attention to the translational dynamics of spherical particles in turbulence [9]. A few considered rotation for neutrally buoyant spheres in turbulence [81, 120, 141]. While these detected rotation, its role on the flow and particle dynamics was not revealed. Also in the freely-rising-sphere-community, rotational effects have been largely neglected [105]. In experiments, often, the particle's mass-density Γ was varied by using various combinations of hollow and solid spheres made from a variety of materials [100, 138, 142, 143]. In doing so, the moment of inertia (MoI or I) varied erratically, introducing variations in the particle's rotational freedom. Therefore, in the present work we ask the question: *Could a change in the moment of inertia of a buoyant particle affect its overall dynamics?* We look into the dynamics of spherical particles in two canonical flow environments: a) in turbulence and b) in quiescent liquid.

5.2 Buoyant particle in a turbulent flow

We begin with the case of buoyant spherical particles released in a turbulent flow. Two spheres were considered, both weighing $\approx 15.2 \pm 0.1$ g, which corresponds to a mass-density ratio $\Gamma \approx 0.88$ in water. The spheres have identical

mass, diameter, and surface properties, which leaves us with identical values for the particle control parameters ($\Gamma \approx 0.88$ and $\text{Ga} \approx 6000$). At the same time, we introduce a subtle modification to the sphere's composition. One sphere was fabricated as a hollow spherical shell, with $I \approx 1.87 \times 10^{-6} \text{ Kg m}^2$, and the other, as a thin spherical shell with a dense metallic core, with $I \approx 1.05 \times 10^{-6} \text{ Kg m}^2$ (see Fig. 5.1(B) & (C)). These result in dimensionless moments of inertia (MoI) $I^* \approx 1.0$ and 0.6 , respectively[†]. The spheres were released in a turbulent flow, which was generated using an active grid placed upstream of the measurement section of the Twente Water Tunnel (Fig. 5.1(A)). The water tunnel was configured to have downward flow in the measurement section, with a Taylor-scale Reynolds number $\text{Re}_\lambda \approx 300$, and the particle size was a fraction of the integral scale of the turbulence ($\Xi \approx 100$). The downward mean flow speed in the measurement volume was adjusted to be comparable to the rise velocity of the particles. This ensured that the buoyant spheres were rising against the flow, enabling us to track long particle trajectories. The particles were imaged using two high-speed cameras placed at a 90 degree angle between them, and recorded at 500 Hz. An analytically prescribed pattern was painted on the spheres, which enabled us to track their orientation in 3D (see Chapter 4). The output of the orientation tracking method, in axis-angle convention, was used to obtain the instantaneous angular velocity and angular acceleration of the spheres. Thus, we obtain a complete description of the particle's dynamics, composed of three translational and three rotational degrees of freedom.

[†]Dimensionless moment of inertia $I^* = I/I_f$, where I_f is the moment of inertia of a fluid sphere with the same volume as the particle.

(C) TNB coordinate system

(A) Tumbling

(B) Fluttering

Figure 5.2: Drawings representative of a typical trajectory traced by the (A) high MoI sphere and the (B) low MoI sphere in a turbulent flow at $Re_\lambda \approx 300$. Zoomed-in views of the rectangular windows pqr s and $jklm$ are also shown. The high MoI sphere tumbles in the flow, wherein the direction of rotation does not change much during motion. This tumbling motion induces a mean horizontal drift for the particle. The low MoI sphere flutters in the flow, wherein the directions of rotation and translation undergo frequent reversals with a frequency comparable to that of vortex-shedding [144]. This fluttering is thought to stabilize its motion to remain vertical in the mean. (C) Sphere motion transformed into a moving TNB coordinate system, where T - the direction of the particle's instantaneous velocity, N - direction of curvature of the particle trajectory, and B - the binormal vector is defined such that $N = (B \times T)$. (D) & (E) Normalized histograms of the angle between the angular velocity vector $\vec{\omega}$ and its time derivative $d\vec{\omega}/dt$ expressed in terms of $\Delta\theta$ and $\Delta\phi$ for the (D) high and low MoI spheres, respectively. Movies S1 & S2 [145, 146] compare the two types of motion.

Figure 5.3: (A) Mean horizontal drift velocities of the high and low MoI spheres in a turbulent flow. (B) Histogram of the angle between the horizontal drifting direction and the vector $\vec{\omega} \times \hat{j}$, where $\hat{j}$ is the unit vector in the vertical direction. The high MoI sphere shows alignment. The low MoI sphere shows two peaks at around 0° and 180° , indicating alignment and anti-alignment. See Appendix C for more details. Movies showing the motion of the two spheres are given in Movies S1 & S2 [145, 146].

The above settings lead to a situation where the four control parameters: Γ , Ga , Ξ , & Re_λ are identical for the spheres we consider. At the same time, the spheres have notably different moments of inertia ($I^* \approx 1.0$ and $I^* \approx 0.6$). This difference leads to differences in the particle's overall dynamics. Evidence for this was first seen while tracking the movement of the spheres. The low MoI sphere undergoes zig-zag motions, while remaining in the middle of the measurement section. In contrast, the high MoI sphere shows a tendency to drift horizontally in the flow and approach the channel walls (see Fig. 5.2(A) & (B) for representative drawings of the trajectories). The mechanism behind these preferential movements becomes clear when we look at the particle's translational and rotational motions simultaneously (see Movie S1 [145] & Movie S2 [146]). The high MoI sphere is in a *tumbling* state, while the low MoI sphere displays a *fluttering* type of motion. In the tumbling state, the directions of rotation and translation do not change much during motion (Fig. 5.2(A)). However, for fluttering, both the rotation and translation undergo periodic reversals of direction (Fig. 5.2(B)), at a rate comparable to the sphere vortex shedding frequency (see Appendix C and Fig. C.1).

Further evidence for the *tumble-flutter* transition can be found by viewing the particle motion from a Lagrangian (TNB) coordinate system, which moves with the sphere (see Fig. 5.2(C)). The orientation of a vector is expressed in terms of the elevation and azimuth angles θ and ϕ , respectively. In Fig. 5.2((D) & (E)), we show the normalized histogram of the angle between the vector $\vec{\omega}$ and its time derivative $\alpha \equiv d\vec{\omega}/dt$, for the two spheres. Here $\vec{\omega}$ and α are the angular velocity and angular acceleration vectors, respectively. An alignment between these vectors would mean the particle's angular velocity

increases in magnitude, but without a change in the direction of rotation. An anti-alignment would mean that the angular velocity decreases, but without a change in the direction of rotation. For any situation where $\vec{\omega}$ changes orientation, the two vectors would not be aligned or anti-aligned. For the high MoI sphere, $\vec{\omega}$ is preferentially aligned with $d\vec{\omega}/dt$ (*single peaked*). $\vec{\omega}$ also aligns with the mean horizontal drifting direction of the particle (see Fig. 5.3(B)), which is evidence that the high MoI sphere statistically tumbles in the direction of its horizontal drift. We believe that this tumbling motion is crucial in establishing the mean horizontal drift for the high MoI sphere (see Fig. 5.3(A)).

The low MoI sphere however shows almost equal probability for $\vec{\omega}$ and $d\vec{\omega}/dt$ to be aligned and anti-aligned (*double peaked* in Fig. 5.2(E)). In addition, $\vec{\omega}$ does not preferentially align with the mean horizontal drifting direction of the particle (see *double peak* in Fig. 5.3(B)). Thus, a fluttering type of motion occurs, which stabilizes the low MoI sphere to remain in the bulk of the water channel flow. An analogy may be drawn to falling disks, strips and paper, where a similar *tumble-flutter* transition has already been observed due to moment of inertia reduction [147, 148]. However, some form of geometrical anisotropy of the particle was considered necessary to induce this transition [149]. The present results suggest that such transitions are possible even for an isotropic body, such as a sphere.

The *tumble-flutter* transition revealed that the high MoI sphere drifts predominantly along a particular direction, while the low MoI sphere (fluttering) undergoes frequent reversals of its direction of motion (Fig. 5.3 and Movies S1 & S2 [145, 146]). The fluid forces responsible for the differences in the dynamics may be gauged from the particle's acceleration statistics. Fig. 5.4(A) shows the probability density function (PDF) of the centripetal acceleration a_N for the two spheres under consideration. A reduction of MoI induces a significant change in the shape of the PDF, with a mean value that is almost double that of the high MoI sphere. When we compare the angular accelerations, the difference is even more dramatic (see inset to Fig. 5.4(A)). If the torques acting on the spheres were comparable, one should expect the angular acceleration to be increased by $\approx 66\%$ ($\mathcal{T} = I d^2\theta/dt^2$). Instead, the *rms* of angular acceleration is increased by $\approx 430\%$. Clearly, particle rotation plays a crucial role in inducing these enhanced accelerations. This is quantified in Fig. 5.4(B), where we show the acceleration a_N conditioned on the magnitude of a rotation-induced lift (or Magnus) force $F_L \propto \vec{v}_p \times \vec{\omega}$ [150]. For the high MoI sphere, a_N appears to be almost independent of the Magnus force. However, a clear dependence is seen for the low MoI sphere. The origin of this becomes clear once we plot the orientation of $\vec{\omega}$ in TNB coordinates (Fig. 5.4(C) & (D)).

Figure 5.4: (A) Probability density function (PDF) of the centripetal acceleration a_N for the high (in black) and low MoI (in red) spheres rising through a turbulent flow at $Re_\lambda \approx 300$. Inset shows the PDF of the angular acceleration α_N for the two spheres. (B) $\langle a_N \rangle_c$, the acceleration conditioned on the magnitude of lift force $(v_p \times \omega)_c$ due to a rotation-induced lift force $C_L |v_p \times \omega|$. The dashed line shows the prediction for the enhancement of $\langle a_N \rangle_c$ due to a rotation-induced lift force $C_L |v_p \times \omega|$. The dotted line shows the nearly constant value of $\langle a_N \rangle_c$ for the high MoI sphere. Normalized histograms showing the orientation of the angular velocity vector in the TNB coordinate system for the (C) high and (D) low MoI spheres in the turbulent flow. The particle rotation aligns strongly with $\vec{B}$ for the low MoI sphere, while it is spread in the NB plane for the high MoI sphere.

For the low MoI sphere, $\vec{\omega}$ aligns along the $\vec{B}$ direction, resulting in $-\vec{v}_p \times \vec{\omega}$ aligning perfectly with the direction of $\mathbf{a}_N$ (Fig. 5.4(D)). Therefore, the acceleration $\mathbf{a}_N$ may be written as $\mathbf{a}_N \approx \mathbf{a}_0 + C_L |\mathbf{v}_p \times \boldsymbol{\omega}|$, as shown by the dashed line in Fig. 5.4(B). For the high MoI sphere, $\vec{\omega}$ lies in the NB plane with no preferential orientation, which explains the absence of correlation in the conditioned acceleration plot (Fig. 5.4(C)). Thus, tuning the moment of inertia has enabled us to alter both the fluid forces and torques acting on a particle, leading to an overall modification of its dynamics in turbulence. However, whether these effects are restricted to a turbulent flow or not is unclear. To address this, we will next look into the motion of buoyant spheres rising through quiescent liquid.

5.3 Freely rising particles

We conducted experiments of spheres rising through quiescent liquid. Fig. 5.5 shows a schematic of the experimental setup used for releasing the spheres (see Appendix C for more details). Fig. 5.6(A) shows the free-rise trajectories and wake-patterns of the high and low MoI spheres rising through quiescent water ($Ga \approx 6000$). At this Ga , the wake is turbulent. The wake is visualized by letting the particle rise through a patch of fluorescent dye. The high MoI sphere (Fig. 5.6(A)(left)) rises with a small amplitude of oscillation and a nearly vertical wake pattern. When I^* is reduced, large amplitude oscillations are triggered, and we witness a spread-out wake behind the particle (Fig. 5.6(A)(right)). A similar behavior is observed at low Ga (see Fig. 5.6(B)). In this case the oscillation amplitude is enhanced, and the wakes differ in the number of vortical structures shed per oscillation. These wake-patterns were previously thought to be formed when the mass-density-ratio Γ

Figure 5.5: Experimental arrangement used for visualizing the trajectories and wakes of freely rising spheres. Fluorescein sodium dye at 5×10^{-3} mol/l concentration is injected at the bottom of the water tank. The sphere is held at the bottom by applying suction.

Figure 5.6: (A) Trajectories and wakes of high (*left*) and low (*right*) moment of inertia spheres at $Ga \approx 6000$. (B) Trajectories and wakes of high (*left*) and low (*right*) moment of inertia spheres at $Ga \approx 500$. In (A) & (B), the wakes were visualized by releasing fluorescent dye. False coloring was adopted based on the intensity of the dye in the wake. Note that the diameter scale ‘D’ is provided in each case. Reducing the moment of inertia at fixed Γ triggers an increase in the oscillation amplitude, along with a change in the wake-pattern. Movie S4 [151] and Movie S5 [152] correspond to the cases shown in (A) and (B), respectively.

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was lowered below a critical value [100]. The enhanced rotational motion due to MoI reduction could be the key factor triggering these modified particle and wake motions.

The above observations have provided some insights into the buoyancy-driven dynamics of freely rising spherical particles. A long-standing debate in the topic relates to the existence a critical mass-density ratio ($\Gamma \approx 0.6$ [106]) marking the onset of large amplitude vibrations for freely rising spheres. While this has been reported in literature, the physical mechanism behind such a sharp transition is puzzling. The main question remains as to how a marginal reduction in the mass-density could give rise to a significant enhancement of the vibration amplitude, since the effective mass of a buoyant particle (i.e. actual + added mass) changes by only a small factor [153]. Our observations suggest that the reason for this could lie in the differences in the particle’s rotational motion, which manifest due to a change in the rotational inertia (since I^* changes along with Γ in most situations).

Similar reportings of a critical mass ratio ($\Gamma = 0.54$ [139]) were made for freely rising (circular) cylindrical particles. Both spheres and circular cylinders are isotropic bodies. Therefore, the insights gained here could be relevant for freely rising circular cylinders as well. In Chapter 6, we will numerically study the trajectories and wakes of freely rising/falling circular cylinders [154].

5.4 Conclusions and outlook

In conclusion, the use of simultaneous 3D particle position and orientation tracking has enabled us to resolve the coupled translational and rotational dynamics of buoyant spheres in a range of flow environments. We have shown that the onset of path-instability can be linked to the rotation of a particle, and that resonance may be triggered (or inhibited) by tuning a particle's moment of inertia. The concept could be exploited in chemical engineering processes where the mixing and transport of nutrients can be effectively enhanced through the introduction of vigorously vibrating particles [128]. Further, the *flutter-tumble* transition we observed for buoyant spheres in turbulence offers a few flow-control opportunities. For instance, particles which migrate toward channel walls could be used to modify the near-wall flow structure [155]. By tuning the moment of inertia, one could design spheres which accumulate near the walls, with potential for drag/heat transfer modifications in dispersed multiphase flow environments.

Dynamics and wakes of freely rising and falling cylinders *

We use direct numerical simulations to study the motion and wake-patterns of freely rising and falling cylinders in quiescent fluid. We show that the amplitude of oscillation and the overall system-dynamics are linked to two parameters: the particle's mass-density relative to the fluid $m^ \equiv \rho_p/\rho_f$ and its relative moment-of-inertia $I^* \equiv I_p/I_f$. Using 144 combinations of m^* and I^* , we comprehensively map out the parameter space covering very heavy ($m^* > 10$) to very buoyant ($m^* < 0.1$) particles. The data collapses into two scaling regimes demarcated by a transitional Strouhal number, $St_t \approx 0.17$. St_t separates a mass-dominated regime from a regime dominated by the particle's moment of inertia. A shift from one regime to the other also marks a gradual transition in the wake-shedding pattern: from the classical 2S (2-Single) vortex mode to a 2P (2-Pairs) mode of wake vortices. Thus, auto-rotation can significantly influence the trajectories and wakes of freely rising isotropic bodies.*

*Based on: V. Mathai, X. Zhu, C. Sun, and D. Lohse, Mass and moment of inertia govern the transition in the dynamics and wakes of freely rising and falling cylinders (under review).

Simulations by Zhu. Analysis and writing by Mathai and Zhu. Supervision by Sun and Lohse, discussion of results and proofreading by everyone.

6.1 Introduction

Path-instabilities are a common observation in the dynamics of buoyant and heavy particles. Common examples are the fluttering of falling leaves and disks, and the cork-screw and spiral trajectories of air-bubbles rising in water [105, 156, 157]. The oscillatory dynamics of such particles can vary a lot depending on the particle's size, shape and its inertia, and the surrounding flow properties. This can be important in a variety of fields ranging from sediment transport and fluidization to multiphase particle- and bubble-column reactors [122–124], since vigorous path oscillations can enhance the drag, and the transport of heat and mass through the fluid.

The simplest case of a rising or falling particle in a fluid is a sphere or a circular cylinder released in quiescent fluid. This problem has traditionally been studied and characterized using two non-dimensional parameters: the solid/fluid density-ratio ($m^* \equiv \rho_p/\rho_f$) and the generalized particle Galileo number, Ga , expressing the ratio between buoyancy and viscous effects [158, 159]. Among these, Ga governs the onset of various kinds of wake-instabilities behind the particle, and m^* governs the motion of the particle in response to these flow instabilities and vortex-induced forces. A number of investigators have studied the influence of these parameters, and various kinds of paths and wakes have been observed for rising and falling bodies [100, 106, 139]. However, despite a fair level of understanding of the mechanisms affecting path- and wake-instabilities, it remains unclear as to what factors precisely trigger vigorous path oscillations for a rising or falling object, such as a sphere or a cylinder.

The motion of a freely rising/falling particle in a fluid is a complex fluid-structure-interaction (FSI) problem. Wake-induced forces cause the particle to move, which in turn changes the flow field around it. This situation holds many similarities to the popular subject of vortex-induced-vibrations (VIV) [108]. In this area, investigators have explored the dynamical response of elastically mounted and tethered bodies in uniform flows. A spring-mass-damper model was often used to make predictions of the oscillatory response of such systems under various conditions [106, 139]. For example, a critical mass-ratio m_{crit}^* was predicted for the sudden appearance of path-oscillations for elastically mounted spheres and cylinders undergoing VIV [25, 108]. For the freely rising body as well, similar predictions were made ($m_{crit}^* = 0.54$ for cylinder, and $m_{crit}^* = 0.61$ for sphere) by modeling it as a spring-mass-damper system with zero spring-stiffness and zero damping [100]. While a dependence on the mass-density ratio is undisputable, others observed wide scatter in their data to the point that there is no consensus with regard to whether a unique critical mass-ratio exists or not [63, 143]. Moreover, from a dynamical point

of view, the existence of a m_{crit}^* lies in contradiction with a fundamental concept: i.e, the motion of a light body through a fluid should be fluid added-mass dominated (since $m_a > m_p$) [22, 61]. Therefore, it remains surprising that a marginal reduction in mass-density alone would trigger the sudden appearance of large-amplitude oscillations, since the effective mass of the system (actual + added mass) is almost unchanged [153]. We will provide a plausible explanation for some of these anomalies.

6.2 Numerical approach

In this chapter, we study the two-dimensional motion of circular cylinders rising or falling through a quiescent fluid. The body geometry (circular cylinder), the state of the surrounding fluid (quiescent), and the imposition of two-dimensionality of the flow make the problem simplified (see Leontini *et al.* [160] for a discussion on the possible three-dimensional effects). Nevertheless, this model problem can provide important clues about the underlying dynamics of buoyancy-driven bodies in general [161]. We performed direct numerical simulations (DNS) using the immersed boundary method. The solver uses the discrete stream-function formulation for the incompressible Navier-Stokes equations [162] with a virtual force implementation [163], which enables us to deal with both light and heavy particles. The code has been extensively validated through prior investigations [164–166]. A rectangular computational domain of the size $100D \times 16D$ is used, with a grid width of $0.01L$ near the cylinder, where D is the cylinder diameter.

The fluid motion governed by the incompressible Navier-Stokes equations may be written in the dimensionless form as:

$$\frac{\partial \mathbf{u}}{\partial t} + \mathbf{u} \cdot \nabla \mathbf{u} = \nabla p + \frac{1}{Re} \nabla^2 \mathbf{u} + \mathbf{f}, \quad (6.1)$$

where $\mathbf{u}$ is the velocity vector, p is the pressure, Re is the Reynolds number, and $\mathbf{f}$ is the Eulerian body-force that is used to mimic the effects of the immersed body on the flow [165, 166]. The direct numerical simulations provide an exact description of the flow-body interaction. For modeling purposes, it is common to adopt a minimalistic description of the fluid flow, wherein the particle motion is written in terms of the forces and moments exerted on the body. The classical Kelvin-Kirchhoff equations expressing linear and angular momentum conservation for the cylinder motion in a fluid may be extended to an incompressible flow containing vorticity [167, 168]. The equations read:

$$(m_p + m_a) \left[\frac{d\mathbf{U}}{dt} + \boldsymbol{\Omega} \times \mathbf{U} \right] = \mathbf{F}_v + (m_p - \rho_f \mathcal{V}_p) \mathbf{g}; \quad (6.2)$$

$$I \frac{d\boldsymbol{\Omega}}{dt} = \boldsymbol{\Gamma}_v; \quad (6.3)$$

where m_p is the particle mass, m_a is the added mass, I is the particle moment of inertia, $\mathbf{U}$ is the particle velocity vector, $\boldsymbol{\Omega}$ is the angular velocity, ρ_f is the fluid density, $\mathcal{V}_p$ is the particle volume, and g is the acceleration due to gravity. $\mathbf{F}_v$ and $\boldsymbol{\Gamma}_v$ represent the vortex forces and moments on the cylinder in a viscous fluid.

Eqs. (6.2) & (6.3) point to two parameter-dependences, namely the particle's mass and its moment-of-inertia. In addition, the rotation rate of the particle Ω is linked to I and couples with eq. (6.2) through a force term. However, the dependence on I was completely neglected for isotropic bodies despite the widespread variation in I in almost all existing experimental studies [59, 100, 106, 139, 140, 143]. We therefore map out the $[m^* - I^*]$ parameter space for rising and falling cylinders.

6.3 Results and discussion

We begin with the case of a very buoyant cylinder ($m^* = 0.1$) rising in a quiescent fluid. Fig. 6.1(a) shows the trajectory of the particle when its rotation is constrained ($\Omega = 0$). The amplitude of oscillation $A/D \approx 0.65$. Next, we let the same particle auto-rotate [169] while it rises through the fluid. Fig. 6.1(b) shows the trajectory for $I^* = 0.1$. We observe an 85% enhancement in the oscillation amplitude due only to relaxing its constraint on rotation. Fig. 6.1(c) & (d) show snapshots of the vorticity field in the vicinity of the particle for the two cases. These observations indicate a clear link between the auto-rotation of the particle and the resulting translational and wake dynamics.

Having seen the effect of rotation on the translational and wake dynamics, we systematically explore the dynamics of heavy and buoyant cylinders, with m^* and I^* both varied in the range $[0.1; 10]$. This spans a wide parameter space covering very heavy ($m^* = 10$) to very light ($m^* = 0.1$) particles. It may be noted that in an actual experiment m^* and I^* are linked by the relation $I^* = m^*(R_{pG}^2/R_{fG}^2)$, which imposes an upper bound on $I^* \rightarrow 2m^*$. Nevertheless, we vary the I^* independently over a wide range $[0.1; 10]$ since it serves as an independent control parameter to modulate the particle's rotational freedom. The resulting translational and rotational amplitudes of vibration are shown in the contour maps in Fig. 6.2(a) & (b), respectively. The translational and rotational amplitudes depend on both m^* and I^* . The change in wake-pattern as we move from upper-right to lower-left part of the contour map in Fig. 6.2(a) & (b) is shown through the sequence (A)–(D) in Fig. 6.2(c). For large m^* , we

Figure 6.1: (a) Oscillatory rise trajectory of a buoyant cylinder ($m^* = 0.1$), with a constraint on rotation ($\Omega = 0$). (b) Oscillatory rise trajectory of the same cylinder when allowed to freely rotate ($m^* = 0.1$, $I^* = 0.1$). Here Ω is the angular velocity, D is the diameter of the particle, and λ is the wavelength of the oscillation. (c) & (d) Vorticity snapshots for the cylinders in (a) & (b), respectively. Angular velocity Ω is zero for the rotationally constrained cylinder. Red-white-blue colormap is used for the vorticity, with blue and red colors indicating the clockwise and counter-clockwise vorticity, respectively. Allowing the cylinder to auto-rotate [169] changes the wake pattern dramatically.

always observe a *2S* (*2-Single*) vortex mode **(A)**, which is identical to that of a cylinder fixed in a uniform flow. However, as we lower the m^* and the I^* , we first observe a *Mixed mode* **((B) & (C))**, which gradually transitions to a clear *2P* (*2-Pairs*) mode of wake vortices **(D)**. From the figure, we see that the *2S-2P* transition can occur for m^* in the range $[0.2 - 1]$, depending on the I^* . While a similar *2S-2P* transition was also observed in the vortex-induced-vibrations literature [170], the transition itself was thought to be sudden at $m^* = 0.54$. Thus, unlike an elastically mounted or tethered particle, the extra rotational freedom present for a freely rising cylinder could be facilitating this gradual transition from *2S* to *2P* vortex mode. The same can be said about the oscillation amplitude, which grows gradually when m^* and I^* are reduced.

Fig. 6.2 demonstrated the existence of regions where the mass did not play a serious role, along with others where the amplitude depended mainly on the mass. However, up until now, there exists no simple explanation for the observed growth of the oscillation amplitude in these regimes. Here we adopt a simplified approach, where we assume that the vortex force on the cylinder (from eq. (6.2)) remains similar in intensity across the cases investigated. As is common in the VIV community, we describe this vortex force as $F_v \approx F_{v0} \sin \omega t$ [22, 144]. Next, we assume that the system dynamics is governed by the vortex shedding frequency ω . For the freely rising case, the

Figure 6.2: Contour map of the (a) translational and (b) rotational amplitudes of motion for buoyant ($m^* < 1$) and heavy cylinders ($m^* > 1$) in the $[m^* - I^*]$ parameter space at $Ga \approx 500$. The shaded regions mark the very high values of I^* , which are unphysical in an actual experiment. This region is primarily of mathematical interest. (c) Vorticity plots for the points marked as (A), (B), (C), and (D) in (a) and (b). A gradual transition in the wake-pattern is observed as we move from upper-right to lower-left regions of (a) and (b). The vorticity is scaled to the $[-1; 1]$ range. Movies showing the different wake patterns are given in the supplemental material [171].

shedding frequency ω can get modified as we change m^* and/or I^* . In the mass dominated regime, ω should primarily be a function of m_p . Thus, the transverse acceleration of the particle is given as $a_p \sim \frac{F_{v0}}{m_p} \sin \omega t$ [22], which leads to a relation for the transverse motion: $x_p \sim \frac{F_{v0}}{m_p \omega^2} \sin \omega t$. Non-dimensionalizing with the particle diameter D and the mean rise velocity U , we obtain the dimensionless amplitude, $A/D \propto 1/(m^* St^2)$, where m^* is the dimensionless mass and $St \equiv \frac{\omega D}{2\pi U}$ is the Strouhal number or equivalently the dimensionless vortex shedding frequency [144]. In Fig. 6.3(a), we plot A/D versus $1/(m^* St^2)$ for the full range of parameters in the present study. The solid blue

curve shows our prediction, where F_v is estimated using the lift coefficient for flow past a fixed cylinder, i.e. $C_L \approx 0.84$. The prediction matches well with the simulation results up to an oscillation amplitude $A/D \approx 0.5$. The inset to Fig. 6.3(a) shows the compensated plot, where the plateau demonstrates the robustness of the scaling.

Figure 6.3: (a) Oscillation amplitude A/D versus $1/(m^* St^2)$ for buoyant and heavy cylinders. The dark blue curve shows the scaling relation $A/D \propto 1/(m^* St^2)$, which holds for $A/D \leq 0.6$. The inset to (a) shows the compensated plot; the plateau for small $1/(m^* St^2)$ demonstrates the robustness of this scaling. (b) A/D versus $1/St^2$ for buoyant and heavy particles. For $A/D > 0.6$, the oscillation amplitude scales as $A/D \propto 1/St^2$ in the compensated plot in the inset to (b). (c) Strouhal number contour map covering the $[m^* - I^*]$ parameter space for the full family of buoyant and heavy cylinders from the present study. A transitional Strouhal number $St_t \approx 0.17$ separates the two scaling regimes shown in (a) and (b). For $St > St_t$, we observe the classical 2S wake mode, which is almost identical to that of uniform flow past a fixed cylinder. For $St < St_t$, the wake transitions from the 2S to a 2P mode. The dashed inclined line in (c) denotes the homogeneous cylinder cases. The shaded region of high I^* values is not possible in an actual experiment. Nevertheless, this region is of mathematical interest, since it shows how at reduced rotations the oscillation amplitude diminishes. Similar scaling regimes were seen at a lower Ga of 220, for which the wake is known to be predominantly two-dimensional.

The prediction for the mass dominated regime is valid up to $A/D \approx 0.5$. Beyond this, changing m^* at fixed I^* has little effect on A/D . At the same time, we observe branching of the oscillation amplitude for different values of the particle moment of inertia (right half of Fig. 6.3(a)). A/D almost doubles in this regime (0.6 to 1.2), solely due to a change in I^* . Therefore, we call this the moment of inertia dominated regime. Here the dependence on m^* is weak, since the fluid added mass outweighs the particle mass. Therefore, the transverse oscillation amplitude in this regime may be expressed without m^* dependence, as $A/D \propto 1/St^2$. A/D versus $1/St^2$ for the full range of parameters is shown in Fig. 6.3(b). A/D scales linearly with $1/St^2$, as demonstrated by the red dashed line. When compensated for this, we observe a plateau for $1/St^2 > 32$ (or equivalently, $St < 0.17$) in the inset. From the plateau, we obtain $C_L \approx 1.57$. The scaling is valid for $A/D \geq 0.6$. Below $A/D = 0.6$, the relation does not hold, since the cylinder mass becomes important. Therefore, one can interpret that the growth in oscillation amplitude in the moment of inertia dominated regime results solely due to the reduction in the dimensionless shedding frequency or St . However, the reduction in St is in turn driven by the enhancement in rotational motion of the cylinder (Fig. 6.1).

The two regimes and their respective behaviors were predicted based on the expectation that the vortex force remains unchanged, but only the Strouhal number depends on m^* and I^* . Thus, a St map (as shown in Fig. 6.3(c)) can summarise the dynamics of the entire family of heavy and light cylinders rising/falling through a quiescent fluid. The clean $2S$ wake mode is found in the mass-dominated regime (*right-side*), and the clean $2P$ mode is restricted to the moment of inertia dominated regime (*lower-left*). The transition from $2S$ mode to $2P$ mode is not sudden, as was found for elastically mounted particles [108, 172], but gradual and marked by the appearance of a *Mixed* wake mode in the $0.5 < A/D < 0.7$ range. The conclusions of the present study hold even at a lower Ga ($Ga = 220$). In this case, however, the ranges of mass- and moment-of-inertia-dominated regimes are slightly different, presumably due to the Ga dependence on the vortex dynamics behind rising cylinders.

Studies on elastically mounted and tethered bodies have proposed that m^* alone determines the oscillation amplitude and wake mode ($2S$ or $2P$). The same models have been extended to freely rising/falling bodies and predicted $m_{crit} \approx 0.54$ (resp. 0.61) for the sudden appearance of vigorous vibration for cylinders (resp. spheres). The present study has revealed that this picture might be an incomplete one for the dynamics and wakes of buoyancy-driven bodies. This is because, unlike an elastically mounted sphere or cylinder, a rising particle has a rotational degree of freedom, which effectively can change its translational dynamics. This perspective might provide an explanation for

Figure 6.4: Vortex-shedding time τ_v^* vs the inverse of the relative speed $1/U_{rel}^*$ for all heavy and buoyant cylinders from the present study. The dashed line suggests a nearly linear dependence. The inset shows the mean rise speed U_{rise}^* and the mean rotational speed $U_{rot}^* \equiv \frac{\langle |\Omega| \rangle D}{2 U_g}$ vs τ_v^* , where Ω is the instantaneous angular velocity of the cylinder, D is the cylinder diameter, and $U_g \equiv \sqrt{gD|1 - \rho_p/\rho_f|}$ is a gravitational velocity scale [22]. Note that $*$ denotes a dimensionless variable, where the non-dimensionalization is done using the length D and velocity U_g .

the scatter reported for the value of m_{crit} in prior experimental studies [106, 139, 143, 173].

The simplified model adopted here has enabled us to understand the motions and wakes of heavy and light cylinders rising/falling through still fluid. From the assumption that the vortex forces remain similar in intensity, but change only in the duration of their action, we can explain the growth in the oscillation amplitude. Since the Strouhal number represents the dimensionless frequency of vortex shedding, the fact that it is reduced means that the vortex shedding is retarded [161]. While the precise mechanism by which rotation induces this is not clear, one can qualitatively explain the behavior. For this, we compare the rotational response of a high moment of inertia cylinder with that of a low moment of inertia cylinder (Fig. 6.1(c) & (d)). The high moment of inertia cylinder resists rotation, while the low moment of inertia cylinder yields to the fluid torques and rotates. On the side where vorticity is shed (see Fig. 6.1(c) & (d)), the particle's rotation is along the mean flow direction. Therefore, the relative speed at the cylinder surface $U_{rel} = (U_{rise} - U_{rot})$ is reduced for the low moment of inertia cylinder, which could cause the developing shear-layer to remain attached for longer. Thus the vortex shedding time $\tau_v \propto 1/U_{rel}$ for the low moment of inertia cylinder, as compared to $\tau_v \propto 1/U_{rise}$ for the non-rotating cylinder. This corresponds to a reduction in the dimensionless frequency for the low moment of inertia cylinder. In Fig. 6.4, we find evidence that the dimensionless vortex shedding

timescale $\tau_v^* \equiv \frac{\tau_v U_g}{D}$ increases almost linearly with the dimensionless inverse relative speed $1/U_{rel}^* \equiv \frac{U_g}{U_{rel}}$, where $U_g = \sqrt{gD|1 - \rho_p/\rho_f|}$ is a gravitational velocity scale [22].

In summary, the present study has demonstrated that the path oscillations of a buoyant particle can be linked to its rotational freedom, as speculated by Ryskin & Leal in 1984 [153]. Our findings remain to be experimentally confirmed using cylinder-rising experiments, such as those conducted in the past [139], but with the rotational inertia as an additional control parameter.

Dynamics of underwater pendulums *

Simple pendulums are used as everyday examples in demonstrations of Newtonian mechanics, wherein the oscillation time period is shown to be independent of the pendulum mass. However, this result is valid only for a pendulum oscillating in vacuum, and when friction is negligible. Here we performed experiments on the motion of heavy and buoyant cylindrical pendulums oscillating under water. We track the oscillation frequency of the pendulums, and its dependence on the pendulum mass. We develop a basic theoretical model, which takes into account the fluid forces and bearing friction. The model predicts the oscillation frequencies and the peak accelerations reasonably well. However, deviations are seen, particularly in the second and subsequent maxima of oscillations. Using Time-Resolved Particle Image Velocimetry (TR-PIV), we reveal that these deviations likely arise due to the disturbed flow created by the pendulum at earlier times. This history effect can lead to an added drag on the cylinder. We incorporate this in the equation of motion, resulting in improvements in the predictions of the second swing and subsequent oscillations.

*Based on: V. Mathai, L. Loeffen, T. Chan, and S. Wildeman, Dynamics of underwater pendulums (to be submitted).

Experiments and analysis by Loeffen and Chan. Supervision by Mathai and Wildeman. Writing by Mathai, Loeffen, Chan, and Wildeman.

7.1 Introduction

The first known study on pendulums dates back to Galileo Galilei in 1605. The story goes that Galileo observed the lamplighter pushing one of the swaying chandeliers in the Pisa cathedral. Galileo timed the swings with his pulse and concluded that, although the amplitude decreased, the time of each swing remained constant. While Galileo's observation might seem trivial in the present day context, it provided the first demonstration of an important model system in Newtonian mechanics.

The equation of motion of a simple pendulum of mass m attached to an infinitely light rigid rod of length L can be written as: $mL \, d^2\theta/dt^2 = -mg \sin\theta$, where θ is the angular displacement from the vertical position, and g is the acceleration due to gravity. Under the small angle approximation, this equation yields the classical result that the period of oscillation $T = 2\pi\sqrt{L/g}$. For finite amplitude oscillations, the time period is larger: $T_\alpha \approx T(1 + \frac{1}{16}\alpha^2 + \frac{11}{3072}\alpha^4)$, where α is the amplitude of oscillation. These demonstrated that the oscillation time period of a simple pendulum is independent of its mass. This elegant result follows from the equivalence of the gravitational and inertial masses in the pendulum equation of motion [174, 175], and is valid when the oscillations occur in vacuum, and in the absence of friction.

While the motion of the pendulum in vacuum is well understood, the typical pendulum dynamics is far from the simple equation of motion presented in textbooks. A few investigators have in recent years looked into the dynamics of heavy objects swinging in viscous fluids. Bolster *et al.* [175] studied a pendulum in water using spherical bobs. They measured the loss in potential energy during the swings to obtain the drag coefficient variation with Reynolds number. This study was conducted in the linear drag regime, where the Stokes drag provided a fair prediction of the forces acting on the spherical bob of the pendulum. In another experiment, Neill *et al.* [174] conducted experiments on steel and brass pendulums oscillating in different fluids. They showed how buoyancy and added mass affect the pendulum motion. Their results agreed fairly well with the motion predicted by taking the buoyancy and added mass into account, implying that viscous corrections played a minor role in the oscillation frequency. A few studies also involved modified pendulums. For instance, Gonzalez and Bol [176] built a physical pendulum with an eddy-current damping system. In this case, a clear frequency shift was detected between the undamped and damped oscillations. In all of the above studies, the researchers focussed on small amplitude vibrations of dense pendulums in viscous fluids. In this regime, the dynamics could be reasonably modeled using the ideal flow approximation, without even including the fluid drag (Stokes drag) term in the equation of motion.

In the present work, we study the large amplitude oscillations of a cylindrical pendulum in water. The experiments are conducted in a regime where the cylinder Reynolds number is of the order of 10^4 . We test the consequences of the non-Stokesian drag and the non-equivalence of gravitational and inertial masses in the equation of pendulum motion. We performed controlled experiments in a water tank to track simultaneously the oscillations and the flow field around a cylindrical pendulum. The control parameter in our experiment is the cylinder to fluid density ratio $m^* \equiv \rho_p/\rho_f$, while the ratio of the cylinder diameter to the pendulum length D/L , the fluid properties (ρ_f and ν), and the release angle ($\theta_0 = 90^\circ$) are fixed. As we will see later, the Galileo number is yet another important parameter governing the ratio of buoyancy to viscous forces. We vary m^* , covering both buoyant ($m^* < 1$) and heavy ($m^* > 1$) pendulums in water. We build on the knowledge of buoyancy, added mass, and drag to propose a basic model for the pendulum motion. The predictions of this model are compared with our experimental observations. We use time-resolved Particle-Image-Velocitometry (PIV) for tracing the flow field around it. Details of the coupled dynamics of the pendulum and the flow around it are presented.

7.2 Experiments

The experiments were conducted using a cylindrical pendulum in water. The length of the pendulum arm L was about 200 mm. The diameter D and the span W of the cylinder used were 32 mm and 300 mm, respectively. This resulted in an aspect ratio $W/D \approx 10$, which could be considered large enough to expect a predominantly two-dimensional mean flow around the pendulum. The cylinder attached to the pendulum arm was constructed with several components: a PVC pipe sealed with two plastic caps, and masses which could be inserted inside the pipe (see Fig. 7.1(c)). This let us study both heavy and buoyant pendulums in water.

The pendulum was placed in a water tank with dimensions 800 mm $\times$ 400 mm $\times$ 500 mm by pivoting on a metal rod of radius $R = 3$ mm, and fixed to the side walls using suction cups. To minimize wall effects, the cylinder always remained at least $1.3D$ from the side walls, $5D$ below the water level for the buoyant case, and $5D$ above the base of the water tank for the heavy case. A 15 minute interval was ensured between successive experimental runs, which ensured that the residual flow from a previous run had damped out. Slightly different setups were used for the heavy and buoyant pendulums. Schematics of the basic experimental arrangement for the heavy and buoyant cases are shown in Fig. 7.1(a) & (b), respectively. While most of the components were the same,

Figure 7.1: Schematics of the experimental setup for (a) heavy and (b) buoyant pendulums. (c) Variation of the cylinder's mass was achieved by inserting masses into cylinder. Note that the rod connecting the cylinder to the axis of the pivot is rigid, and the cylinder itself did not rotate independent of the rod. This meant that the moment of inertia of the cylinder had no significant influence on the pendulum dynamics. (d) Experimental setup used for performing PIV measurements.

the main difference was in the release mechanism. The release mechanism consisted of a semicircular cavity with two curved metal wire segments. The cylinder was held against the semicircular cavity and locked in place by the metal wires. With this arrangement, it was possible to release the pendulum by pulling the wire. This ensured a gentle release of the pendulum from above the water level without creating disturbances in the water. By inverting these parts and using a long hook, the buoyant pendulums could also be released without disturbing the water. The release position was at 90° angle from the vertical. To record the position of the cylinder during the oscillations, a high-speed camera (Photron Fastcam 1024PCI) was used. The cylinder movement was detected in the images using a circle detection method [22].

In addition to tracking the pendulum motion, we performed Time-Resolved Particle Image Velocimetry (TR-PIV) to measure the velocity field in the fluid surrounding the cylinder. Fluorescent tracer particles $\approx 50 \mu\text{m}$ in diameter were released in the flow, and a high speed Laser (Litron LDY-303HE) with appropriate cylindrical lens was used to generate a sheet of light passing

Figure 7.2: Force diagram on the cylindrical pendulum at an instant when it swings to the right. The pendulum arm is rigid. As shown, the cylindrical pendulum bob experiences buoyancy $\vec{F}_B$, weight $\vec{F}_g$, drag $\vec{F}_D$, and a normal force due to the pendulum arm $\vec{F}_N$; while at the pivot, the pendulum experienced some bearing friction $\vec{F}_f$.

through the mid-section of the cylinder. Mirrors were placed as shown in the Fig. 7.1(d). The two mirrors ensured that the shadows cast by the cylinder were removed by reflected light. This enabled us to have nearly uniform illumination all around the cylinder. A double-frame camera (Imager sCMOS) at 25 fps and inter-pulse duration of about 2 ms was used to obtain the time-resolved velocity field. The experimental setup used for PIV measurements is shown in Fig. 7.1(d).

7.3 Equation of motion

The equation of motion of a pendulum oscillating in a fluid can be expressed by Newton's law in the angular form:

$$I \frac{d^2\theta}{dt^2} = \sum \tau, \quad (7.1)$$

Here, θ represents the angular position of the pendulum with respect to the equilibrium position, I is the moment of inertia of the system, and $\vec{\tau}$ represents the torque acting on the system (see Fig. 7.2). The pendulum mass is mostly concentrated at a distance L from the axis of rotation. Even for the lightest case studied here, the moment of inertia of the rod about the pivot axis contributed to less than 6% of the total moment of inertia of the system. Therefore, the system could be modeled as a point mass with a moment of inertia: $I_p = m_{eff}L^2$, where m_{eff} is the effective mass of the system. When

the pendulum accelerates through the fluid, it must also accelerate some fluid surrounding it. This influence can be modeled using the added mass concept. The effective mass of the system m_{eff} is the sum of the pendulum mass m and the fluid added mass m_a ; $m_{eff} = m + m_a$.

The torques acting on the system can be written as $\sum \vec{\tau} = \sum \vec{F} \times \vec{r}$, where $\vec{F}$ are the individual forces, and $\vec{r}$ the corresponding arm lengths. The forces acting on the system are gravity $F_g \equiv \rho_p \mathcal{V} g$, buoyancy $F_B \equiv \rho_f \mathcal{V} g$, fluid drag F_D , and a bearing friction $F_f \approx \mu_f F_N$. Here, ρ_p and ρ_f are the particle and fluid densities, respectively, $\mathcal{V}$ is the cylinder volume per unit length, μ_f is the friction coefficient, and $F_N \approx F_g - F_B$ is the normal force acting at the bearing. The Galileo number $Ga \equiv \sqrt{g D^3 |m^* - 1|} / \nu$ gives an estimate of the Reynolds number built on some characteristic length scale D of the body and on the gravitational velocity scale $v_g \equiv \sqrt{g D |m^* - 1|}$. We estimated Ga of the cylinder (see Table 7.1). This suggests that the pendulum motion will be in the high Reynolds number limit. Therefore the drag force may be written as $F_D = \frac{1}{2} C_D \rho_f A_p v_p^2$, where C_D is the drag coefficient, A_p is the projected area of the cylinder, and $v_p \equiv L d\theta/dt$ is the instantaneous cylinder velocity. The equation of motion after non-dimensionalization with respect to $m^* \equiv \rho_p / \rho_f$ and $\tilde{t} \equiv t \sqrt{g/L}$ can be written as

$$m_{eff}^* \frac{d^2 \theta}{d\tilde{t}^2} = -k \sin \theta - c \left| \frac{d\theta}{d\tilde{t}} \right| \frac{d\theta}{d\tilde{t}} - h \frac{d\theta}{d\tilde{t}} \left| \frac{d\theta}{d\tilde{t}} \right|^{-1} \cos \theta, \quad (7.2)$$

where the constants, m_{eff}^* , k , c and h are defined as:

$$\begin{aligned} m_{eff}^* &= (m^* + m_a^*) \\ k &= (m^* - 1) \\ c &= \frac{1}{2} C_D \frac{A_p L}{\mathcal{V}} \\ h &= \mu_f (m^* - 1) \frac{R}{L} \end{aligned}$$

In order to predict the motion of the pendulum, m_a^* , C_D , and μ_f have to be set. In our experiments, the ratio of the length of the cylinder to its diameter D was nearly 10. Hence, the ideal flow added mass of an infinite cylinder, $m_a^* = 1.0$, was chosen to begin with. The drag coefficient to expect from our experiments could be estimated using the range of values of Ga for the cylinders (see table 7.1). As a first approximation, Ga could be viewed as being comparable to the cylinder Reynolds number. Ga lies in the range [3800 – 35800]. Based on the C_D vs Re curve for the flow past a circular cylinder, we choose $C_D = 1.2$ [177]. Since the contact between the moving

Figure 7.3: (a) An image of a heavy pendulum a few moments after release. (b) Angular position (θ) vs time for two cases with different mass ratios. The oscillation frequencies and amplitude decay rates are notably different. (p1) and (p2) approximately correspond to the maximum deceleration points for the two cases.

Table 7.1: Ga and Re for some of the tested m^* . Here, $\text{Re} = U_0 D / \nu$, where U_0 is the maximum speed attained by the pendulum calculated from the model based on $C_D = 1.2$, and neglecting the added mass and friction. Note that the calculated values of Re are comparable to Ga.

	Heavy								Buoyant			
m^*	4.98	4.15	3.65	2.82	2.20	1.79	1.37		0.95	0.75	0.54	0.33
$\text{Ga} \times 10^{-3}$	35.8	31.8	29.2	24.2	19.6	15.8	10.9		3.8	9.0	12.2	14.6
$\text{Re} \times 10^{-3}$	33.7	30.9	28.8	24.7	20.6	16.9	11.9		4.2	10.0	13.7	16.6

and the stationary parts of the pendulum at the pivot was steel on steel, μ_f was expected to lie in the range $[0.2 - 0.3]$ for lubricated steel-on-steel contact. We used $\mu_f = 0.2$, based on a few tests performed for one of the heavier pendulums in air, since in this case the fluid drag and added mass effects were at a minimum.

With the parameters chosen, the predictions of the model were examined. We first looked into the maximum cylinder Reynolds number predicted by the model. As indicated in table 7.1, Re is comparable to Ga, indicating that the approach is self consistent. This also suggests that the pendulum attained a speed comparable to the terminal speed of a falling cylinder in 2D, i.e $v_g \simeq \sqrt{gD |m^* - 1|}$. Next, we focus on the full motion of the pendulum. Fig. 7.3(b) compares the predicted motion for heavy pendulums with $m^* = 1.16$ and $m^* = 4.98$. The amplitude decay, and the frequency both depend on the mass of the cylinder chosen. This is rather expected for an underwater pendulum, since the mass-density ratio m^* appears in the dimensionless equation of motion eq. (7.2). With this in mind, we explore the predictions for the range

Figure 7.4: (a) Contour-plot of θ against t and m^* . Note the asymmetry of the graph about $m^* = 1$. (b) Phase portraits comparing the angular velocity ω_{pend} vs the angular position θ for a buoyant pendulum ($m^* = 0$) with that of a heavy pendulum ($m^* = 2$) having identical driving force $|F_B - F_g|$. (c) Phase portraits of five heavy cases. With increasing m^* , both the maximum angular position and maximum angular speed increase. (d) Surface plot of the figure shown in (a). The heavier the pendulum, faster the decay of its oscillation amplitude.

of m^* covering buoyant and heavy pendulums.

7

7.4 Results and Discussion

Fig. 7.4(a) shows a contour plot of the oscillation angle against time for continuous variation of m^* . We notice a clear asymmetry about $m^* = 1$ line, which separates the heavy cases from the buoyant cases. This suggests that the dynamics are very different for comparable driving force ($|\vec{F}_B - \vec{F}_g|$) on the heavy and buoyant sides. For instance, the oscillation frequency is higher for a buoyant cylinder as compared to a heavy cylinder with identical driving ($|m^* - 1|$). This can be rationalized by taking just the buoyancy and added mass into account, since $f \propto \sqrt{\frac{|m^* - 1|}{m^* + 1}}$ even for the small amplitude oscillations of a simple pendulum moving through in ideal (frictionless) flow [174]. However, many

details of the motion are missing in this prediction. In Fig. 7.4(b), we show the phase portraits of two cases, one buoyant ($m^* = 0$) and other heavy ($m^* = 2.0$), with identical $|m^* - 1|$ and hence the same $|\vec{F}_B - \vec{F}_g|$. The buoyant pendulum with ($m^* = 0$) accelerates quickly, achieving a high angular speed in a short period of time. At the same time, owing to its low inertia, which in this case comes entirely from the fluid (or m_a^*), the motion damps out quickly. Therefore, the peak angular displacement at first minimum θ_0 is lower for the buoyant pendulum as compared to the heavy pendulum ($m^* = 2.0$), while the peak angular velocity achieved ω_{peak} is higher for the buoyant cylinder.

Fig. 7.4(c) shows the phase portraits of the angular velocity vs angular displacement for the heavy cases. The larger the m^* , the greater are the initial acceleration, the maximum angular speed achieved, and the number of oscillations to getting damped. However, this plot does not feature the time evolution of the oscillations. The duration and decay of the oscillations for the heavy cases are evident from Fig. 7.4(d). Interestingly, the heavier the pendulum, the quicker the damping of the oscillations, and this occurs despite their higher inertia. The interplay between the higher speeds achieved by the heavier pendulums (as seen in Fig. 7.4(c)), and the non-linear growth of the drag in proportion to the square of the pendulum speed is responsible for this. However, in the upper limit of $m^* \gg 1$, the damping rate should reduce with m^* , as commonly seen for pendulums oscillating in air. This implies the existence of an optimal heavy pendulum ($m_{opt}^* > 1$), for which one observes the quickest damping.

Fig. 7.4 showed some key differences in the dynamics of heavy and buoyant pendulums as predicted by the model equation of motion. However, a one to one comparison with experiment is necessary to establish the validity of the predictions. In Fig. 7.5(a), we present the normalized frequency of oscillation from the experiments along with the predictions of the model (eq. 7.2). As expected, the frequency of oscillation has a clear dependence on the mass ratio m^* . When the pendulum mass density is close the cylinder density (or $m^* \rightarrow 1$), the frequencies are low. For heavy cylinders, the frequency increases with m^* , and asymptotically approaches the frequency of a large amplitude simple pendulum, given by the horizontal dash-dotted line in Fig. 7.5(a). The predictions of the model (eq. 7.2) are given by the blue and red dashed curves. The model provides a fair prediction of the oscillation frequency for both heavy and buoyant cylinders. Here we have used a constant drag coefficient $C_D \approx 1.2$ from literature, which is based on the mean drag coefficient for a cylinder at the Reynolds number in our experiments. The blue curve corresponds to the prediction when the potential flow added mass $m_a^* = 1.0$ is used. This slightly underpredicts the oscillation frequency. A better fit to the data is obtained

Figure 7.5: (a) Comparison of the normalized frequency of oscillation f^* from the experiments and the predictions of the model. Here the normalization is done with the frequency of a simple pendulum for small amplitude oscillations, i.e $f_{sp} = \frac{1}{2\pi} \sqrt{\frac{g}{L}}$. An added mass coefficient $m_a^* = 0.53$ provides the best agreement between model and experiment. Note that for undamped oscillations ($C_D = 0$, $\mu_f = 0$), the model underpredicts the oscillation frequency. (b) Peak pendulum acceleration, a_{max} plotted against m^* for the buoyant and heavy cases in our study.

for $m_a^* = 0.53$. Thus, for $C_D = 1.2$, $m_a^* = 0.53$, and $\mu_f = 0.2$, the predictions of the model are in reasonable agreement with the experiments for both buoyant and heavy cylinders.

It is interesting to compare the predictions of the model (Eq. 7.2) when the drag and the bearing friction are ignored altogether. The dashed green curve shows the frequency prediction when the added mass alone is accounted for in the model. Neill *et al.* [174] had previously observed that for a heavy pendulum undergoing small amplitude oscillations, such a model should suffice. However, in the present situation, the model grossly underpredicts the oscillation frequency of heavy pendulums (see Fig. 7.5(a)). Clearly, the non-linear drag in the present situation brings in major changes to the oscillation frequency, in contrast to prior observations of heavy pendulums undergoing small amplitude oscillations in viscous fluids.

The frequency of oscillation represents an averaged quantity for the pendulum motion considered here. A more sensitive quantity would be the peak acceleration a_{max} of the oscillating pendulum. a_{max} also signifies the instant when a maximum net force acts on the pendulum. We estimated the peak acceleration using the position vs time data for the cylinder. Following the initial instants of release of the cylinder, a peak in the acceleration is experienced when the cylinder reaches the end of the first swing. This occurred near the points marked as (p1) & (p2) in Fig. 7.3(b). In Fig. 7.5(b), we compare the experimentally measured peak acceleration with the predictions of

the model (eq. 7.2) for two different added mass coefficients, $m_a^* = 1.0$ and $m_a^* = 0.53$, respectively. The peak acceleration changes significantly with m^* , and the model predictions are in reasonable agreement with the experimental measurements. Interestingly, the value of added mass has only a minor role on a_{max} . This suggests that the fluid drag has a leading role in the peak deceleration of the pendulum.

With the adequacy of the model in the form of Eq. (7.2) established, we plot the time evolution of angular position (θ) for one of the heavier pendulums with $m^* = 4.98$ (see Fig. 7.6(a)). For comparison, we provide the predictions of the model for $C_D = 1.2$, $m_a^* = 0.53$, and $\mu_f = 0.2$. As expected, the oscillation frequency matches between the experiment and the model. The angular position predictions of the model agree with the experiments up to the end of the first swing (point marked as (q) in Fig. 7.6(a)). This is also consistent with our earlier observation of the peak acceleration (shown in Fig. 7.5(b)), which occurred just prior to the end of the first swing. However, deviations are seen in the second peak of oscillation, marked as (r) in Fig. 7.6(b). The deviation persists for all of the subsequent oscillations.

Fig. 7.6(a) indicates a situation when the first swing of the pendulum is captured fairly well by our model, while a deviation is seen for the second swing and subsequent oscillations. This was also the case for other mass ratios studied. Thus the model is unable to capture some important physical effects present in our experiments. To understand the possible origin of these deviations, we employ PIV to quantify the flow around the cylinder during its motion. The normalized horizontal velocity, and the normalized vorticity plots shown in Fig. 7.6(c), (d), and (e) correspond to the time instants marked as (p) , (q) , and (r) , respectively, in Fig. 7.6(a).

Fig. 7.6(c) shows the velocity and vorticity fields at a time instant when the cylinder is moving towards the right to its first minimum. The wake behind the cylinder is clearly visible, while the flow ahead of the cylinder has negligible vorticity or is nearly irrotational (Fig. 7.6(c)(right)). The wake is unsteady and is expected to induce unsteady forces on the cylinder during its swing. However, one can expect that the mean drag on the cylinder at this instant is fairly predicted by the non-linear drag term F_D in the equation of motion (eq. 7.2). Therefore, the first swing of the pendulum is captured well by our model. Next we focus on a later instant in time (point marked as (q) in Fig. 7.6(a) & (b)), when the cylinder just completed its first swing and is returning to the left. Fig. 7.6(d) shows the velocity and vorticity fields at this instant. In this case, however, the cylinder is moving towards a disturbed background flow. The horizontal velocity plot (see *left-side* plot of Fig. 7.6(d)) indicates that the cylinder faces an incoming flow to the right. The effect of this disturbed flow is not taken into account in our model. Therefore, as the cylinder travels further through the disturbed flow, it may experience a greater resistance than what is predicted by the drag term in our model. Fig. 7.6(e) shows the flow field at the instant when the cylinder completed its reverse swing (point marked as (r) in Fig. 7.6(a) & (b)). At this instant, the original wake in front of the cylinder appears to have dissipated.

Figure 7.6: (a) Graph of angular position θ vs time t for a heavy pendulum with $m^* = 4.98$. The red curve represents the model prediction based on Eq. (7.2) for $C_D = 1.2$, $m_a^* = 0.53$, and $\mu_f = 0.2$. (b) shows the second swing in greater detail. The model over-predicts the maximum amplitude (point marked as (r) in (a) & (b)). (c), (d) and (e) show the normalized horizontal velocity U_x/v_0 (left) and normalised vorticity $\omega/(v_0/D)$ (right) at the three instants marked in subfigure (a) as (p), (q), and (r), respectively. Here, v_0 is the measured maximum speed of the pendulum during its first oscillation. U_x pointing to the right and ω pointing out of the page are positive.

Figure 7.7: Correction factor $E(t)$ vs θ/θ_0 for the return swing of the first oscillation. The legends show the different m^* cases. $E(t) > 1$ initially, so the wake's effect on the pendulum's motion is expected to be significant. Inset shows the percentage error in the first maxima predicted by the original model, and those due to a revised model, which takes a representative flow velocity U_f also into account. The predictions are improved after including this correction.

One of the approximations in using the drag force in Eq. (7.2) was that the cylinder velocity nearly equals the relative velocity v_{rel} between the cylinder and the flow. However, as we saw from the PIV snapshots, the flow develops a mean velocity after the first swing of the pendulum. This disturbed flow in the wake could result in v_{rel} deviating significantly from v_p , which might be responsible for the overprediction of the maxima of second and subsequent oscillations by the model. We begin by estimating the possible deviations in the drag term due to the incoming flow velocity. We estimate a correction factor E , which accounts for the possible error due to the usage of v_p , instead of v_{rel} . Here, $v_{rel} = (\vec{U}_f - \vec{v}_p)$, where $\vec{U}_f$ is defined as the mean flow velocity in a rectangular window with dimensions $D \times D$, with its centre at a distance of $1.5D$ ahead of the cylinder. In Fig. 7.7, we plot E vs θ/θ_0 for the return swing of the pendulum for various mass ratios, where θ_0 is the peak amplitude of the return swing. While the choice of the window size was rather arbitrary, it provides a reasonable estimate of the error in the drag term. During the initial phase of the return swing, E exceeds unity for all the cases, suggesting that the effect of the disturbed wake on the pendulum's motion is significant. With time, E decreases, and the effects due to the disturbed flow become negligible. The inset to Fig. 7.7 quantifies the percentage error in the peak

amplitude prediction for pendulums of different m^* . The two cases shown are the predictions of the original model, and those of a corrected model, which uses v_{rel} in the drag term. The error in the first maxima of oscillation is decreased due to the inclusion of the correction in the drag term.

7.5 Conclusions

In this work, we have studied the large amplitude oscillations of heavy and buoyant cylindrical pendulums in water. The oscillation frequency and the peak acceleration of the pendulum depend on the pendulum mass density ratio m^* . Accounting for buoyancy and added mass alone, as done in previous investigations [174], does not provide a fair prediction of the dynamics. Therefore, we developed a basic theoretical model, which accounts for the non-linear drag force and the bearing friction, in addition to the buoyancy and the added mass. We used the potential flow added mass coefficient $m_a^* = 1.0$, and the steady flow drag coefficient C_D from literature [177]. With the inclusion of these terms, the model predictions are in reasonable agreement with the experimental observations for a wide range of mass ratios. The predictions are further improved for a reduced $m_a^* = 0.53$, which provides the best fit to our experimental data. The reduced added mass coefficient as compared to the ideal flow situation is logical, since the flow is viscous. Moreover, in choosing the m_a , we made the assumption of two-dimensionality of the flow. In reality, the pendulum is of finite length, and the cylinder Reynolds numbers are high enough that the flow in the wake is surely turbulent, i.e three-dimensional.

While the oscillation frequency and peak acceleration are captured well, the model overpredicts the peak oscillation amplitudes for the second and subsequent oscillations. We used particle image velocimetry (PIV) to quantify the flow around the cylinder during its motion. The PIV flow field reveals the disturbed flow surrounding the cylinder, the effect of which is not accounted in the model. We included a correction into our model based on the mean speed of the flow in the vicinity of the cylinder. This changes the drag force term to include the instantaneous relative velocity between the cylinder and the flow: $v_{rel} = \vec{U}_f - \vec{v}_p$, instead of the cylinder velocity $\vec{v}_p$ used in the original model. The predictions of the peak oscillation amplitude are improved with the inclusion of this correction term.

The theoretical model presented in the chapter is basic. Yet it provides a fair prediction of the pendulum dynamics for a range of mass density ratios covering heavy and buoyant pendulums oscillating in water. Nevertheless, there are several limitations to this approach. For example, the value of the drag coefficient $C_D = 1.2$ is for the steady flow past a circular cylinder at a

high Re . In reality, the cylinder Reynolds number varies in time. Also, the unsteady forces due to vortex shedding behind the cylinder are not included in our model. Another source of uncertainty comes from the modeling of the bearing friction. The bearing friction could also have a dependence on the speed of rotation. Further, the PIV based corrections improved the predictions of the model. However, this required input from the experimentally measured flow field, which can be difficult to obtain. Fully resolved direct numerical simulations (Immersed boundary [89], or Physalis [84, 88] etc) could provide better insights into the flow-induced forces. On a different note, the present experiments have shown the importance of fluid drag for a bluff body oscillating in a fluid. An interesting extension to the study would be to use a hydrofoil shape, which could reduce the contribution from drag. In this case, we expect the frequency and dynamics may be captured well by a simpler model, which takes only the buoyancy and added mass into account.

Liquid agitation by bubbles rising within an incident turbulent flow *

This work reports an experimental characterisation of the flow properties in a homogeneous bubble swarm rising at high-Reynolds-numbers within an incident turbulent flow. We vary the gas volume fraction α and the velocity fluctuations u'_0 of the carrier flow, respectively, in the ranges $0\% < \alpha < 0.9\%$ and $2.3 \text{ cm/s} < u'_0 < 5.5 \text{ cm/s}$. The so-called bubble parameter ($b \equiv \frac{V_r^2 \alpha}{u_0'^2}$) is used to compare the ratio of the kinetic energy of the bubbles to that of the incident turbulence, and is varied from 0 to 1.3. Here V_r is the bubble rise velocity. The liquid velocity fluctuations increase with b , for $b > 0.27$. For $b < 0.27$, we observe turbulence attenuation. Conditional measurements of the vertical velocity downstream of the bubbles allow us to disentangle three regions that have specific statistical properties, namely the primary wake, the secondary wake, and the far field. While the fluctuations in the primary wake are similar to those of a single bubble rising in a liquid at rest, the statistics of the velocity fluctuations in the far field follow Gaussian distribution, similar to the one produced by homogenous and isotropic turbulence at the largest scales. In the secondary wake region, the conditional probability density function (pdf)

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Experiments and analysis by Almeras and Mathai. Writing by Almeras and Mathai. Supervision by Lohse and Sun. Discussion of results and proofreading by everyone.

of the velocity fluctuations is asymmetric and shows an exponential tail for the positive fluctuations and appears to be nearly Gaussian for the negative fluctuations. The overall agitation thus results from the combination of these three contributions and depends mainly on the bubble parameter.

8.1 Introduction

Numerous industrial processes, such as the Fischer-Tropsch and reaction catalysis make use of the agitation induced by bubbles to enhance mixing and mass transfer. Thanks to their buoyancy, bubbles rise at a high velocity through the liquid, thereby inducing perturbations in the liquid. However, the turbulence induced by bubbles, commonly called pseudo-turbulence or ‘burbulence’, differs a lot from homogeneous-isotropic-turbulence (HIT) or wall-bounded turbulence. In the last two decades, several studies have been conducted on this subject, both experimentally [19, 30, 178–182] and numerically [47, 181, 183–186], leading to a clearer understanding of the bubble-induced liquid agitation. The liquid phase fluctuations due to pseudo-turbulence result mainly from two contributions, one from the bubble wakes themselves, and the other, from the non-linear interactions between the bubble wakes [187, 188]. However, these two contributions do not have the same role. While the fluctuations contained in the wake are mostly upward and cause a positive skewness of the pdf of the velocity fluctuations in the vertical direction, most of the liquid agitation comes from the interaction between the wakes, which results in exponential tails of the velocity pdfs [179]. A consequence of this interaction between the wakes is that the velocity fluctuations have a non-linear dependence on the gas volume fraction, as: $V_r \alpha^{0.5}$, where V_r is the relative rising velocity of the bubbles with respect to the carrier fluid, and α is the gas volume fraction. Another feature of pseudo-turbulence is that it is inherently anisotropic, meaning that the vertical fluctuations produced by the bubbles are stronger as compared to the horizontal ones.

While pseudo-turbulence has been adequately characterized, in most natural and industrial settings, bubble-induced agitation occurs in the presence of surrounding turbulence. This leads to a more complex agitation, since it results from the coupling of two sources of turbulence. A very different resulting agitation can be observed in such a bubble swarm rising in a turbulent background flow, depending on the nature of the surrounding turbulence, on the ratio of the energies produced by the two sources, and on their characteristic length- and time-scales. In this work, we focus on the agitation produced by a bubble swarm rising within a nearly homogenous and isotropic turbulent flow. One parameter, which is commonly used to characterise such a turbulent flow, is the so-called *bubblance* parameter b . It is defined to be proportional to $\frac{V_r^2 \alpha}{u_0'^2}$, where $u_0'^2$ is the variance of the velocity fluctuations produced by the external turbulence in the absence of bubbles [19, 112, 182, 189]. By definition, $b = 0$ corresponds to single-phase flow, whereas $b = \infty$ corresponds to a bubble swarm rising in a quiescent liquid (pseudo-turbulence). A wide range

of bubble parameters are thus possible between these two extreme limits, where both pseudo-turbulence and homogeneous and isotropic turbulence are coupled. However, not many studies exist in this domain apart from the spectral analyses of Refs. [182] & [19].

One important feature of turbulent flows is the wide range of length and time scales. The energy contained at these scales can be represented in the spectrum of velocity fluctuations. For homogeneous-isotropic-turbulence, the spectrum of velocity fluctuations shows the classical $-5/3$ scaling for the inertial sub-range in both wavenumber and frequency domains. However, for pseudo-turbulence the spectrum of the velocity fluctuations displays a signature -3 scaling [182, 187]. This originates when the fluctuations induced by the neighbouring bubble wakes interact and get dissipated. Consequently, the -3 scaling is believed to start at a wavenumber $\lambda^{-1} = \frac{C_d}{d}$, where C_d is the drag coefficient of the rising bubble, and d is its diameter [179, 181].

The purpose of the present work is to study the properties of the liquid agitation in turbulent bubbly flows, in a regime where both the incident turbulence (nearly HIT) and the pseudo-turbulence play a role, and to disentangle the different sources of agitation. This leads us to explore the changes to the wake behind the bubble when the bubble parameter is varied. A similar range of bubble parameter was studied by Prakash *et al.* [19]. However, they focussed only on the normalized velocity PDFs and energy spectra. Here, in addition to these analyses, we present a conditional analysis of the statistical properties. This allows us to distinguish three regions in the flow (the primary wake, the secondary wake, and the far field), whose physical properties are different.

The chapter is organised as follows. First, the experimental setup and the measurement techniques are presented in section 8.2. Then, in section 8.3 we discuss the operating conditions in order to characterise the incident turbulence (without bubbles), and the properties of the bubbles. The dynamics of the liquid phase is then explored in two stages. The first stage consists in investigating the overall properties of the liquid phase (section 8.4). In the second stage, we explore the conditional statistics of the flow (section 8.5). A discussion is carried out to link the overall and conditional statistics (section 8.6) before concluding with the main findings and future recommendations (section 8.7).

Figure 8.1: (a) The Twente Water Tunnel (TWT) facility, where the experiments were performed. (b) A zoomed-in view shows the measurement section with the experimental arrangement used in the present study.

8.2 Experimental set-up and instrumentation

8.2.1 Experimental setup

The experiments were performed in the Twente Water Tunnel, which is a vertical water tunnel designed to study two-phase turbulent flows. A schematic of the setup is shown in Fig. 8.1 (a) [20, 21, 59]. The measurement section is 2 m high, 450×450 mm in cross-section, and made of three glass walls allowing optical access. An upward mean flow passes through an active grid which was positioned below the measurement section (Fig. 8.1 (b)). The active grid generates nearly homogeneous and isotropic turbulence in the measurement section. By varying the rotation speed of the active grid and the liquid mean flow, different turbulence intensities were achieved. Bubbles were injected below the active grid by means of 9 islands, each containing 69 capillary tubes of inner diameter 0.12 mm. The islands were placed in a regular 3×3 matrix arrangement, and at a distance of 3 m below the measurement section. This presumably gives enough time for the rising bubbles to distribute nearly homogeneously in the cross-section of measurement. The gas volume fraction is varied from 0.25% to 0.93% by changing the gas flow rate. An ascending turbulent bubbly flow was thus produced.

8.2.2 Instrumentation for the gas-phase characterisation

The gas phase characterisation consists in measuring the global gas volume fraction, the bubble diameter, and the bubble rising velocity. For that purpose, different measurement techniques have been used. The global gas volume fraction α was measured by using an inverted U-tube [112]. We also observed that the global gas volume fraction was comparable to the one measured locally by an optical fiber probe, suggesting that the bubble swarm was homogeneously distributed. The diameter and the rising velocity of the bubbles were measured by means of an in-house dual optical fiber probe. The optical fiber probe detects the passage of the bubble interface at two measurement points, which are separated vertically by a distance, $\delta = 3.41$ mm. The bubble velocity was then calculated as $V_r = \delta/\Delta t$, where Δt is the time interval of the passage of the bubble interface between the two probes. The bubble diameter was estimated based on the time during which the leading probe remained in the gas phase. More details about the signal processing used can be found in Colombet *et al.* [190].

In order to get more details about the gas-phase behaviour, in particular about the aspect ratio of the bubbles, imaging was performed by using a high-speed camera (Photron FASTCAM SAX2), with a 105 mm macro lens focused at the center of the measurement section. The recordings were made at an acquisition frequency of 750 fps and with an exposure time of 1/14035 s. The aperture of the camera was fully open, resulting in a narrow depth of field. The depth of field was estimated to be around 1 cm by using a tilted calibration plate. An image-processing method was developed to measure the diameter and the velocity of the rising bubbles. The bubbles were detected by thresholding based on the spatial intensity gradient of the raw images. In order to detect only single bubbles, two criteria have been imposed on the detected elements: (a) on the equivalent diameter $d_{eq} = \sqrt{4A/\pi}$, where A is the area of the detected element, and (b) on the solidity. The solidity criterion is given by A/A_{cv} , where A is the area of the region suspected to be the bubble, and A_c is the area of the smallest convex polygon engulfing the region of the suspected bubble. Here we used a solidity threshold 0.95, which is comparable to the value used in literature [191].

8.2.3 Instrumentation for the liquid phase characterisation

The liquid phase has been characterised by means of Constant Temperature Anemometry (CTA) for both single-phase and two-phase flow measurements. The CTA probe was placed in the middle of the cross-section of measurement, and the signal was acquired at a frequency of 10 KHz for at least 25 min.

U	u'_0	Re_λ	τ_η	T_L
m/s	cm/s	-	s	s
0.278	2.3	177	0.083	3.6
0.272	3.1	242	0.066	3.9
0.270	3.3	262	0.061	3.9
0.265	3.4	265	0.059	3.8
0.466	3.5	216	0.044	2.3
0.460	4.6	315	0.037	2.9
0.458	5.1	342	0.034	2.8
0.461	5.5	361	0.030	2.6

Table 8.1: Summary of the flow parameters for single-phase measurements. Here, U is the mean flow velocity, u'_0 is the standard deviation of velocity fluctuations, Re_λ is the Taylor Reynolds number, and τ_η and T_L are the dissipative and integral length scales of the turbulent flow, respectively.

It is well known that using CTA in a bubbly flow induces disturbances on the velocity measurement when the bubbles collide with the probe, generating spurious peaks in the velocity signal. Different methods have been adopted to remove the influence of these spurious peaks and to measure only the liquid-phase fluctuations [19, 30, 112, 192]. In the present study, bubble collisions are detected by thresholding the time derivative of the raw velocity signal [193]. In fact, the time derivative of the velocity signal during a bubble collision is typically well over 150 m/s^2 . Therefore, a threshold value of 150 m/s^2 was used. This value is well above any liquid phase velocities encountered in the water tunnel flow and therefore, must occur due to the bubble collisions. Thanks to this operation, the liquid phase is distinguished from the bubble collisions. The mean liquid velocity U , the standard deviation of the velocity fluctuations u_{rms} , and the probability density function (*pdf*) are thus calculated by considering the velocity in the liquid phase only. Before calculating the spectrum of the velocity fluctuations, the gas-phase areas of the signal are removed and replaced by linear interpolation. More details about this operation may be found in Refs. [194–196].

8.3 Operating conditions

8.3.1 Turbulence characterisation

Before investigating turbulent bubbly flows, we will describe the properties of the incident turbulence generated by the active grid. In this study, eight turbulence intensities are investigated, corresponding to two mean flows ($U = 0.27 \text{ m/s}$ and $U = 0.46 \text{ m/s}$) and four rotation speeds of the active grid (see Ta-

Figure 8.2: Characterisation of the homogeneous isotropic turbulence for different operating conditions (single phase flow). (a) PDF of the velocity fluctuations normalised by the standard deviation. (b) Normalised spectrum of the velocity fluctuations.

ble 8.1). Liquid velocity fluctuations are thus increased from $u'_0 = 2.3$ cm/s to $u'_0 = 5.5$ cm/s by increasing the mean flow and/or the rotation speed, resulting in a variation of the Taylor-Reynolds number Re_λ from 177 to 361. For the eight cases, the turbulent velocity fluctuations are nearly homogeneous and isotropic, as also seen from previous studies using the same setup [19, 22, 30]. In fact, once normalised by the corresponding standard deviation, the pdf of the velocity fluctuations shows a nearly Gaussian distribution (Fig. 8.2 (a)). The spectrum of the velocity fluctuations presents a $-5/3$ scaling covering about two decades (Fig. 8.2 (b)) within the frequency range $[1/T_L; 1/\tau_\eta]$ (see Table 8.1), where T_L is the integral time scale calculated from the dissipation rate ϵ ($T_L = (\frac{3}{4}C_0 \frac{\epsilon}{k})^{-1}$, with $k = \frac{3}{2} u'^2_0$, $C_0 = 2.1$). The energy dissipation rate ϵ is estimated from the second order longitudinal structure function, which is calculated by means of the Taylor hypothesis, similar to the methods followed in Refs. [21–23, 35]. The dissipative time scale $\tau_\eta \equiv (\nu/\epsilon)^{1/2}$ is finally estimated for each turbulence level (table 8.1). It must be noted that the above procedures involve Taylor’s frozen flow hypothesis and the assumption that the second order structure function $D_{LL} = C_2(\epsilon r)^{2/3}$ [6]. Typical errors in these estimates lie within 25%, and have been quantified through prior investigations [20, 21].

Figure 8.3: Gas phase characterisation for different turbulence intensities. (a) Mean bubble diameter d vs gas volume fraction α for different turbulence intensities u'_0 . (b) Bubble size distribution for different turbulence levels at the lowest gas volume fraction in our experiments. (c) Aspect ratio χ of bubbles vs mean bubble diameter d . (d) Relative bubble velocity as a function of the gas volume fraction α for different turbulence intensities. The circle symbols denote the measurements made from the dual optical probe, and the triangle symbols denote estimates obtained from high-speed imaging. The color code reflects the rms of liquid velocity fluctuations u'_0 of the incident turbulence.

8.3.2 Bubble properties

For the eight turbulence levels considered in this study, the gas volume fraction is varied from 0% to 0.9 % for $U = 0.27$ m/s, and from 0% to 0.6 % for $U = 0.46$ m/s. The mean bubble diameter ranges from 2.0 mm to 3.6 mm and depends both on the gas volume fraction and the turbulence level (see Fig. 8.3(a)). For all levels of turbulence, the mean bubble diameter increases with the gas volume fraction. Moreover, a strong influence of the rotation speed of the active grid on the bubble diameter can be observed: the faster the rotation speed, the smaller are the bubbles. This trend has already been observed in prior studies [19, 35]. The size distributions of the bubble diameter are presented in figure 8.3 (b) for each level of turbulence at the lowest gas

volume fraction investigated. The standard deviation of the bubble diameter is about 25% – 30% of the mean diameter for all the gas volume fractions and turbulence levels studied. Concerning the aspect ratio of the bubbles, it ranges from 1.6 to 2.0, depending mainly on the rotation speed of the grid and weakly on the gas volume fraction. As shown in Fig. 8.3 (c) the aspect ratio of the bubble is fully controlled by the bubble diameter: smaller bubbles are more spherical compared to larger ones. The Weber number, $We = \frac{\rho V_r^2 d}{\sigma}$ in the given range varies from 1.8 (nearly spherical) to 7.1 (considerably deformed).

The relative bubble rise velocity V_r measured using the optical fiber is shown in Fig. 8.3 (d) as a function of the gas volume fraction for the eight turbulence levels studied. The relative velocity measured by the dual optical probe was comparable to the one evaluated from the high-speed images (not shown here). The variation between the two techniques was within $\pm 5\%$ for the moderate volume fractions studied here. The rising velocity is calculated as the difference between the measured bubble velocity V_b and the liquid mean flow velocity U , i.e $V_r = V_b - U$. Under the present operating conditions, the relative rise velocity ranges from 0.27 m/s to 0.39 m/s and depends both on the turbulence level and the gas volume fraction. The effective bubble Reynolds number based on Fig. 8.3(a) & (d) lies in the range 600 – 900. For any chosen turbulence intensity, the relative rise velocity decreases with the gas volume fraction, as has already been observed for pseudo-turbulence [179, 190]. More experiments would be required to have a deeper understanding of the dependence of V_r on the gas volume fraction and the turbulence level. This falls outside the scope of the present study. Nevertheless, we measured the relative rise velocity for the present operating conditions, since this is a key parameter affecting the liquid agitation.

8.4 Dynamics of the liquid phase within the swarm

8.4.1 Standard deviation of the velocity fluctuations

Fig. 8.4 shows the standard deviation of the velocity fluctuations u_{rms} normalised by the incident turbulent fluctuation u'_0 as a function of the gas volume fraction α for different levels of turbulence. We notice that even after normalisation by the fluctuation of the incident turbulence, the standard deviation of the velocity fluctuation still shows a strong dependence on both α and on u'_0 . For $u'_0 < 3.5$ cm/s, the velocity fluctuations are larger than the one produced by the incident turbulence, and increase with α whereas they are lower for $u'_0 > 4.6$ cm/s. These two behaviours can be explained by the bubble parameter b , which compares the energy of the fluctuations produced by the bubble swarm, which scales as αV_r^2 , to the energy of the fluctuations pro-

Figure 8.4: Standard deviation of the liquid velocity fluctuations in the vertical direction normalised by u'_0 as a function of the gas volume fraction for different levels of turbulence.

duced by the incident turbulence $u'_0{}^2$. The bubble parameter can thus be expressed as:

$$b \equiv \frac{V_r^2 \alpha}{u'_0{}^2}. \quad (8.1)$$

Note that $b = 0$ corresponds to single phase flow, whereas $b = \infty$ is for pseudo-turbulence cases.

Figure 8.5 shows the normalized intensity of velocity fluctuations u_{rms} vs b . A strong correlation exists between the two. Regardless of the level of turbulence, all the data collapse, showing a non-monotonic evolution with the bubble parameter. For $b < 0.27$, the velocity fluctuations are smaller than the one produced by the incident turbulence. Thus, adding a small number of bubbles in a very strong incident turbulent flow leads to a reduction in the intensity of the velocity fluctuations, or turbulence attenuation [197, 198]. However, for $b > 0.27$, an enhancement of the velocity fluctuations compared to single phase flow is observed. The lack of data in the range $0 < b < 0.13$ limits us from having a clear understanding of the transition between a single-phase flow and a turbulent bubbly flow. It is however difficult to investigate this range of b , since it requires smaller gas volume fractions. The current capillary islands in the water tunnel could not produce a homogeneous bubble swarm at such low gas flow rates. We will therefore focus only on the turbulent bubbly flows for $b > 0.13$. In this regime, the normalised velocity fluctuations increase monotonically with the bubble parameter. Two regimes can be observed, separated by a critical bubble parameter $b_c \approx 0.7$ (see the inset to Fig. 8.5). For $b < b_c$, the normalised velocity fluctuations evolve roughly as $u_{rms}/u'_0 \propto b^{0.4}$, whereas they increase much faster for $b > b_c$, approximately as

Figure 8.5: Standard deviation of the liquid velocity fluctuations in the vertical direction normalised by u_0' as a function of the bubbance parameter. The errorbars denote the minimum and maximum values of five equally sampled datasets. Inset shows the same plot on log-log scale.

$\propto b^{1.3}$. This change in the evolution of u_{rms}/u_0' could be due to the increased contribution of the bubble wakes when the bubbance parameter exceeds b_c . The reason for proposing the existence of a critical bubbance parameter will be explained in section 8.5.

In figure 8.5, an interesting transition from turbulence attenuation ($u_{rms}/u_0' < 1$) to turbulence enhancement ($u_{rms}/u_0' > 1$) occurs at $b = 0.27$. While the mechanism behind this is unclear, the value of b at which this occurs might not be a coincidence. To rationalise this, we resort to an alternate definition of the bubbance parameter, based on the actual level of fluctuations in bubbly flow to that of the incident turbulence. As measured by Riboux *et al.* [179], the actual level of fluctuation in a homogeneous bubbly flow $u_b'^2 \approx \gamma^2 V_r^2 \alpha$, where $\gamma = 1.94$ is a prefactor. If one were to use this definition, the transition from turbulence attenuation to turbulence enhancement occurs at bubbance ~ 1 . However, one may note that the prefactor γ could change depending on the flow condition. Therefore, we chose to use the bubbance parameter without γ^2 for simplicity of definition (Equation 8.1). Regardless of the definition chosen, the prefactor has little effect on the main conclusions of the present study.

8.4.2 PDF of the velocity fluctuations

We will now discuss how the structure of the velocity fluctuations depend on the bubbance parameter b . For that purpose, the pdf of the vertical velocity fluctuations normalised by the standard deviation u_{rms} is plotted in Fig. 8.6 for

Figure 8.6: PDF of the liquid velocity fluctuations normalised by u_{rms} for different bubbance parameters. The skewness $\mathcal{S}$ for increasing b is shown as inset.

b in the range 0 – 1.3. For $b = 0$, which corresponds to single phase flow, the pdf of the normalised velocity fluctuations is expected to be nearly Gaussian. For turbulent bubbly flows ($b > 0$), a skewness develops with stronger fluctuations for positive values, but the distributions are nearly Gaussian for the negative and more probable values. Interestingly, the pdfs do not exhibit exponential tails, as is the case for pseudo-turbulence ($b = \infty$). It has been shown in Ref. [179] that the exponential tails are mainly due to the agitation produced by the wakes and their interaction. In a turbulent bubbly flow, the agitation does not result solely from the interactions of the wakes, since the incident turbulence plays a role as well.

The positive skewness of the pdf is also observed experimentally for different operating conditions and is known to be the signature of the bubble wakes [19, 179]. However, in Ref. [19], they could not observe a clear trend with respect to the b parameter. Here we demonstrate that by increasing the bubbance parameter, the positive skewness increases to a maximum value for $b \approx 0.7$ (inset to Fig. 8.6). This can be interpreted as a stronger contribution of the bubble wakes as the bubbance parameter increases. Beyond $b = 0.7$, $\mathcal{S}$ appears to decrease a bit before saturating at large values of the bubbance parameter.

Figure 8.7: Energy spectra of the velocity fluctuations for (a) $b = 0.13$ and (b) $b = 1.3$. The black curves in (a) and (b) denote the spectrum for the single-phase case before the addition of bubbles.

Figure 8.8: Ratio of the integral time scale (λ/V_r) imposed by the bubbles over the dissipative time scale (τ_η) in the flow as a function of the bubble parameter b .

8.4.3 Frequency spectrum of the velocity fluctuations

We now investigate the frequency spectrum S_{zt} of the velocity fluctuations for different values of b . For each operating condition, the spectrum is calculated by averaging at least 180 spectra, each spanning a recording time

Figure 8.9: Energy spectrum of the velocity fluctuations normalised by f_c for different values of the bubbance parameter b .

duration of 8 s, and normalised by u_{rms}^2 . First, the two-phase flow spectrum is compared to the corresponding single phase one, for $b = 0.13$ (Fig. 8.7(a)) and $b = 1.3$ (Fig. 8.7(b)). As already seen in subsection 8.3.1, the classical $-5/3$ decrease can be observed for single phase flow in the range of frequency $[1/T_L; 1/\tau_\eta]$. For the turbulent bubbly flows, the $-5/3$ scaling is still observed for the frequency range $[1/T_L; 1/\tau_\eta]$, and is then followed by -3 scaling for higher frequencies. As is known from Refs. [179, 182], the -3 slope is the signature of the bubble-induced turbulence and develops for length scales smaller than $\lambda = d/C_{d0}$, where C_{d0} is the drag coefficient of a single rising bubble. The characteristic time scale T_{pseudo} imposed by the bubble swarm may then be calculated as $T_{pseudo} = \lambda/V_r$, and the corresponding frequency scales are shown by the green lines in Fig. 8.7(a) & (b). Irrespective of the bubbance parameter, we observe that the -3 scaling subrange starts at the frequency $1/T_{pseudo} = V_r/\lambda$. Since the frequency V_r/λ imposed by the bubbles is larger than the Kolmogorov frequency $1/\tau_\eta$ for the range of bubbance parameters studied here (Fig. 8.8), we can observe a distinct separation of the time scales in the flow, i.e the incident turbulence acting for the time scales larger than τ_η , and the bubble-induced turbulence for time scales lower than λ/V_r , which results in two distinct slopes. However, for $b = 1.3$, the spectrum is deformed in the $-5/3$ subrange and presents a bump at $f_c \equiv 0.14V_r/d \approx 12$ Hz (Fig. 8.7(b)). Fig. 8.9 presents the spectrum of the velocity fluctuations normalised by f_c for different bubbance parameters. It appears that the bump around f_c is present only for $b > 0.27$. This excitation frequency has already been observed in Ref. [184] for pseudo-turbulence. The authors attributed it to the injection

of energy by the collective wake instability. However, this frequency is also comparable to the vortex shedding frequency $f_v \equiv St V_r/d$ of a single bubble at $Re_b \approx 500$, where the Strouhal number $St \approx 0.13$ [199, 200]. It is important to also note that this frequency is close to the one defined by Prakash *et al.* [19] as $f_b = \frac{V_r}{2\pi d}$. We note that the time-scale separation in the present experiments is too narrow to reveal the true origin of this frequency, but this is an interesting issue for future investigation.

8.5 Conditional properties of the liquid phase

8.5.1 Velocity disturbance behind the bubbles

We will now study the velocity disturbance behind individual bubbles. The conditional velocity u_c is calculated by averaging the velocity disturbance just behind each detected bubble (at least 1087 bubbles were detected for each measurement point). It is important to mention that the neighbouring bubbles could induce perturbations on the conditional velocity. The upstream disturbance induced by the next bubble is excluded by shortening the detected sample $3\frac{d}{V_b}$ s before the following bubble is detected. The three diameters distance was chosen based on the analysis of Roig *et al.* [201], wherein the influence of the upcoming bubble was seen to be minor. Here also, we checked that the perturbation was negligible at an upstream distance of three bubble diameters. Finally, Taylor hypothesis is used to convert the temporal signal to a spatial one by using the mean rising velocity of the bubbles V_b [4]. The velocity disturbance behind the bubbles normalised by the turbulent fluctuations u'_0 is shown in Fig. 8.10 for different bubble parameters. It is clear that the bubble parameter strongly affects the structure of the wake. Both the length of the wake and its shape depend on the bubble parameter. For low bubble parameter ($b < 0.7$), the wake is around $5d$ long and presents a single exponential decrease. However, when the bubble parameter is larger than 0.7, the wake length considerably increases, reaching almost $15d$ to $20d$ (see inset to Fig. 8.10). Moreover, the conditional mean velocity presents an exponential decrease with a change in slope at a critical distance z_c at which $u_c(z_c) = u'_0$. We can thus distinguish two domains in the wake: a primary wake for $z < z_c$, referred to as PW, and a secondary wake (SW) for $z > z_c$.

A primary wake (PW) and a secondary wake (SW) were also observed in Ref. [202] for a single bubble in a turbulent pipe flow. The PW was defined as the region which showed a fast velocity decrease, and the transition from PW to SW occurred when the velocity deficit was of the order of the incident turbulence. Following a similar approach, we quantify the PW and the SW

Figure 8.10: Velocity disturbance behind a bubble in the turbulent bubbly flow. Inset shows the velocity disturbance normalised by $(u_c(0) - U)$, which reveals that the primary wake region is not affected by b .

Figure 8.11: Classification of the bubble-wake into primary wake (PW) and secondary wake (SW) regions. Here $k = 1/\ln(10)$ is a prefactor in the slope shown in the above plot.

in our case by defining two characteristic length scales each: a characteristic length l and a characteristic decay length L (Fig. 8.11). In the primary wake ($0 < z < l_{pw}$), the velocity disturbance is well described by an exponential decrease $u_c(z) - U = (u_c(0) - U) \times \exp(-\frac{z}{L_{pw}})$, where L_{pw} is the characteristic decay length of the primary wake. The characteristic length, l_{pw} , is calculated by using the exponential fitting as $u_c(z = l_{pw}) = u'_0$. Concerning the secondary

wake ($l_{pw} < z < l_{sw}$), the velocity disturbance follows also an exponential decrease, which can be expressed as $u_c(z) - U = (u_c(l_{pw}) - U) \times \exp(-\frac{z-l_{pw}}{L_{sw}})$, where L_{sw} is the characteristic decay length of the secondary wake, and l_{sw} is the length of the secondary wake. The length of the secondary wake l_{sw} is defined from the exponential fitting for $u_c(z) - U = 0.05 u'_0$. The threshold was chosen as high as possible, while ensuring that the far-field was nearly Gaussian. In this regard, the value $0.05 u'_0$ seemed to be a reasonable choice. For the fitting of the secondary wake (see figure 8.11), we checked manually for all the cases that the fits were reasonable.

We now focus on the evolution of the characteristic lengths with the bubble parameter (Fig. 8.12 & 8.13). Concerning the primary wake (PW), the characteristic length l_{pw} increases roughly linearly with b , while the characteristic decay length L_{pw} is almost independent of b (Fig. 8.12 (a) & (b)). The characteristic length l_{pw} and the decay length L_{pw} are thus controlled by two different mechanisms. The characteristic length of the primary wake is controlled by the ratio between the two sources of turbulence and thus by b . On the contrary, since the decay length L_{pw} does not depend on b , it means that the velocity deficit behind a bubble is not affected by the incident turbulence and is only controlled by the perturbation induced by the bubble. This is further clarified by the inset to Fig. 8.10, which shows that the velocity deficit in the primary wake is not affected by the level of external turbulence. These findings thus extend the observations of Legendre *et al.* [202] for a single bubble in turbulence to the case of a dilute bubble swarm. Concerning the secondary wake, it presents different characteristics when b increases (Fig. 8.13 (a) & (b)). In particular, the length of the secondary wake l_{sw} is almost constant for $b < 0.7$, after which it suddenly increases, reaching $\approx 20 d$ for $b > 0.8$. The decay length of the secondary wake L_{sw} presents a similar behaviour. It is constant and close to the decay length of the primary wake for $b < 0.7$, after which it suddenly increases. Consequently, the development of the secondary wake is discontinuous, with a threshold at around $b = 0.7$, suggesting that the turbulence induced by bubbles should be strong compared to the external turbulence for the bubbles to develop a large secondary wake. However, the underlying physical mechanisms triggering the secondary remain unclear.

8

8.5.2 Conditional pdf

We now carry out a specific study in the vicinity of the bubble. Three regions are distinguished: the primary wake, ($0 < z < l_{pw}$), the secondary wake ($l_{pw} < z < l_{sw}$), and the far field ($l_{sw} < z$), referred to FF later. Conditional statistical properties of the velocity fluctuations, such as the standard deviation and the pdf, are calculated in each region by retaining the velocity

Figure 8.12: Characteristics of the primary wake. (a) Characteristic length. (b) Characteristic decay length.

Figure 8.13: Characteristics of the secondary wake. (a) Characteristic length. (b) Characteristic decay length.

fluctuations detected in those regions specifically. Fig. 8.14 (a)–(d) presents the conditional pdf normalised by the corresponding standard deviation of the velocity fluctuations u_{rms} for each region and different bubble parameters. We observe that the shape of the conditional pdf strongly depends on the region under consideration. In fact, the conditional pdfs of the primary wake (PW) are asymmetric, presenting a strong skewness for positive values due to the strong perturbation induced by the bubbles. Remarkably, the pdf of the

Figure 8.14: PDF of the velocity fluctuations for different regions of the wake behind the bubble. (a) & (b) Primary wake (PW), (c) secondary wake (SW), and (d) far-field (FF).

velocity fluctuations in the PW, once normalised by V_r , is similar to the one for a single bubble rising in a quiescent liquid (see Fig. 8.14 (b) & [178]). This means that the primary wake disturbance arises mainly from the bubble and is not much modified by the external turbulence. Concerning the conditioned pdf of the secondary wake, they also are asymmetric, since the tails are exponential for positive fluctuations and Gaussian for negative ones (Fig. 8.14 (c)). The secondary wake is thus believed to result from the interaction of the wake with the external turbulence, generating exponential tails for the conditional

Figure 8.15: Relative characteristic volume fraction (β) of the primary wake (β_{pw}), secondary wake (β_{sw}), and far field (β_{ff}). The dashed, dotted, and dash-dotted curves give our interpretation of the evolution of the PW, SW, and FF, respectively. The fractions sum to 1, i.e. $\beta_{pw} + \beta_{sw} + \beta_{ff} = 1$.

pdf for the positive fluctuations only. In contrast, the conditional pdfs in the far field are Gaussian (Fig. 8.14 (d)).

8.6 Discussion

The overall agitation in a turbulent bubbly flow results from a combination of the contributions from the primary wake, the secondary wake and the far field, weighted by their respective volumes. Each conditional pdf presents different properties as shown in Fig. 8.14. The question now is how to evaluate the volume fraction of each region, especially since we only measured the length of the wakes in the vertical direction, and the lateral spread of the wake is unknown. In order to estimate the volume of each region, the typical distance between two bubbles d_{2b} is introduced. Assuming that the bubble are oblate ellipsoids with an aspect ratio χ and homogeneously distributed in space, the mean distance between two bubbles can be estimated as $d_{2b} = \frac{2d}{3\chi^{2/3}\alpha}$. The non-dimensional volume fraction β of the primary wake (resp. secondary wake) may be written as $\beta_{pw} = \frac{l_{pw}}{d_{2b}}$ (resp. $\beta_{sw} = \frac{l_{sw}}{d_{2b}}$), and the volume fraction of the far field as $\beta_{ff} = 1 - (\beta_{sw} + \beta_{pw})$. Fig. 8.15 presents the volume fraction of each region as a function of the bubble parameter b . We can observe that even though the weight of the primary wake increases with b , it constitutes only a small fraction of the total liquid volume, less than 10%, for the full

Figure 8.16: Superposition of the conditional PDF weighted by their respective volumes for (a) $b = 0.13$, (b) $b = 0.27$, (c) $b = 0.74$, and (d) $b = 1.3$.

range of b investigated. In contrast, the secondary wake occupies almost 50% of the total volume when b is larger than 0.7. The secondary wake thus plays a significant role at larger bubble parameters.

Fig. 8.16 presents the superposition of each conditional pdf weighted by their respective volumes for four contrasted bubble parameters: $b = 0.13$, 0.27 , 0.74 , and 1.3 . We observe that superposing the three contributions allows us to get a realistic estimate of the total liquid agitation in the turbulent bubbly flow, even if the spreading of the wakes has not been taken into account. Thus, for low b , the agitation comes mainly from the far field, whereas it comes from both the far field and the secondary wake at larger values of b . However, irrespective of the value of b , the primary wake contributes mainly to the large positive fluctuations ranging from $4u_{rms}$ to $8u_{rms}$, causing the skewness of the

pdf of the liquid fluctuations. This is similar to the observations of Ref. [178] for pseudo-turbulence.

8.7 Conclusions

The hydrodynamic properties of a turbulent bubbly flow have been studied experimentally by varying the level of the homogeneous and isotropic turbulence produced by the active grid (from $u'_0 = 2.3$ cm/s to 5.5 cm/s) and the gas volume fraction (α from 0% to 0.93%). We found that the bubblance parameter is a suitable parameter to characterise such turbulent bubbly flows, as it compares the level of liquid disturbance induced by the bubbles to that due to the external turbulence in absence of bubbles. The bubblance parameter is thus defined as $b \equiv \frac{V_r^2 \alpha}{u_0'^2}$, and varied in the range 0 – 1.3.

We used the Constant Temperature Anemometry (CTA) technique to look into the overall hydrodynamic properties of the turbulent bubbly flow. The spikes in the signal, which correspond to the gas-phase, were detected and removed, and we studied the evolution of the liquid phase velocity fluctuations with b . When normalised by the incident turbulent fluctuations, the velocity fluctuations collapse onto one curve, but show a non-monotonic evolution with b . For $b < 0.27$, an attenuation of the turbulence is observed, while for $b > 0.27$, the standard deviation of the velocity fluctuations are enhanced. Interestingly, this transition from turbulence attenuation to turbulence enhancement occurs when the rms of the bubble-induced fluctuation equals the rms of the liquid fluctuations of the incident turbulence. For the range $0.13 < b < 1.3$, we identified two regimes, with a transition appearing at a critical bubblance parameter, $b_c = 0.7$. For $b < b_c$, the normalised velocity fluctuations u'/u'_0 evolve as $\propto b^{0.4}$. For $b > b_c$, the velocity fluctuations show a steeper increase with b , i.e $\propto b^{1.3}$. The pdf of the velocity fluctuations, once normalised by the standard deviation, follows a Gaussian distribution for negative and more probable values, but with a slight skewness for positive values. This skewness increases with b .

To reveal the origin of the two regimes and the skewness of the pdfs, we performed a conditional analysis on the turbulent bubbly flow. For this, we looked into the velocity fluctuations developing behind individual rising bubbles in the flow. This allowed us to distinguish three regions within the liquid phase: a primary wake, a secondary wake, and a far field. The primary wake, which is the region just behind the bubble, is mainly controlled by the bubble perturbation. As a consequence, the characteristic decay length of the primary wake is constant across all the bubblance parameters investigated. Furthermore, the pdf of the conditional velocity fluctuations of the primary wake normalised by

the relative bubble velocity is very similar to that of a single bubble rising in quiescent liquid. In contrast to this, the conditional pdf in the far field presents a Gaussian behaviour, meaning that it is mainly controlled by the incident turbulence. The secondary wake region results mainly from the interaction of the primary wake with the surrounding turbulence. The conditional pdf of the velocity fluctuations provides an illustration of this interaction, as they have nearly Gaussian tails for the negative fluctuations and exponential tails for the positive ones. The properties of the secondary wake, such as its characteristic length and characteristic decay length present a sharp transition at the critical bubble parameter $b_c = 0.7$. For $b < 0.7$, the length of the secondary wake is $\approx 5d$, and the remainder of the distance to the next bubble is occupied by the far-field, which has Gaussian fluctuations. This results in the weaker dependence of the overall velocity fluctuations on b , corresponding to the $u_{\text{rms}} \propto b^{0.4}$ regime. For $b > 0.7$, the secondary wake is more developed and its length suddenly increases, reaching almost $20d$. This means that the contribution of secondary wakes starts to become important for $b > 0.7$, and therefore we see a stronger b dependence in this regime ($u_{\text{rms}} \propto b^{1.3}$).

The conditional pdfs also provide important clues about the origin of the skewness of the pdf. It turns out that the skewness is due to the primary wake region behind the bubbles, since most of the fluctuations in the primary wake lie in the range of 4 to 8 u_{rms} . Furthermore, we observed a stronger skewness with increasing b , mainly due to the increase of the relative length of the primary wake with b . While the skewness of the pdf is the signature of the primary wake, the Gaussian shape of the pdf for negative fluctuations results from the far field and the secondary wakes fluctuations.

Another quantity that can help to disentangle the effect of the incident turbulence from the one induced by bubbles is the spectrum of the velocity fluctuations. It can also provide some crucial information on the time scales at which the incident turbulence and the one induced by bubbles act. In fact, for the present experimental conditions, the lower frequencies range: $1/T_L$ to $1/\tau_\eta$ shows a $-5/3$ scaling, which is the expected slope of the spectrum for homogeneous and isotropic turbulence. For frequencies larger than V_r/λ , where $\lambda = d/C_d$, the spectrum presents a -3 scaling, which is the signature of the turbulence induced by bubbles. Normalising the spectra with a frequency comparable to the collective wake-instability of the bubble swarm ($f_c \approx 0.14V_r/d$) makes the two ranges more clearly distinguishable for all the bubble parameters studied here. This frequency is also comparable to the vortex shedding frequency of a single bubble. The true origin of this is still unclear.

The present approach, which is based on the decomposition of the overall fluctuations into three contributions, has allowed us to understand the physical

origin of dependences of the velocity fluctuation and the skewness of the pdf on b . In this regard, the sudden growth in the length and the decay length of the secondary wake at $b = 0.7$ is crucial to interpret our observations. As we have found here, the development of the secondary wake significantly enhances the overall turbulence in the liquid phase. Therefore, a better understanding of the physical mechanisms which trigger the secondary wake is required. This could be a key parameter for industrial spargers and bubble-column reactors, where more efficient mixing is desirable. In future work, we intend to link the liquid agitation behaviour observed here to the mixing of a passive scalar in turbulent bubbly flows.

Conclusions and outlook

In this thesis we have studied the dynamics of buoyant particles in a variety of flow environments using a combination of experiments, numerics, and theory. We focussed on two classes of particles: (a) small buoyant particles in turbulence, and (b) finite-sized buoyant particles (and bubbles) in turbulence. The dynamics of the former can be captured by modeling them as passively advected particles in turbulent flow. For finite-sized buoyant particles (and bubbles), the flow modification by the particle significantly affects their dynamics.

Conclusions

In Chapter 2, we have investigated the possibility of using small Stokes number ($St \ll 1$) microbubbles ($\approx 150 \mu\text{m}$) as trackers of turbulent acceleration. These bubbles were comparable in size to the smallest length scale of the flow, and had a response time much shorter than the Kolmogorov time scale of the flow. We used a moving camera setup to measure the Lagrangian statistics of acceleration of the bubbles in a turbulent flow generated in the Twente Water Tunnel facility. Microbubbles show much higher accelerations as compared to similar-sized neutrally buoyant tracer particles. They also show quicker decorrelation time for acceleration. While these results appear similar to those found previously for finite St bubbles in turbulence [24], they cannot be explained as an inertia effect, since the bubbles in our experiment have a very small St . We have used one-way coupled Euler-Lagrangian direct numerical simulations (DNS) to compare with the experimental observations. The DNS provides insight that the enhanced accelerations occur due to the finite drift velocities of these microbubbles in the presence of gravity. This can also explain the reduction in correlation time of acceleration. We showed that the Froude number, if defined as a_η/g [14], cannot capture these gravity induced effects. Instead, $Fr \equiv \frac{a_\eta}{(1-\beta)g}$, where the density of the particle is included through β . This enables us to quantify the acceleration variance and correlation time for arbitrary density particles at small Stokes numbers. Further, we have shown that the mean vertical drift of these particles through the flow can result in

a reduction in the intermittency of acceleration as compared to ideal tracers. Thus, a tiny bubble or droplet is not necessarily a good tracer of turbulent acceleration. This can be important for bubbles and droplets that drift through the turbulent oceans ($g/a_\eta \approx 100 - 1000$) and clouds ($g/a_\eta \approx 10 - 100$) [70]. On the practical side, the findings point to an important consideration when choosing bubbles or droplets for flow visualization and particle tracking in turbulent flows [43].

The next step was to study the effects of finite-size for buoyant particles in turbulence. To this end, in Chapter 3, we studied the Lagrangian dynamics of rigid spherical particles with a density lower than that of water. These provided the first experimental measurements in a previously unexplored regime of buoyant particles with density ratio $\Gamma \equiv \rho_p/\rho_f$ in the range $0.02 - 0.92$, and size-ratio $\Xi \equiv d_p/\eta$ in the range $4 - 50$. We have used a novel experimental strategy to obtain long trajectories of these buoyant particles. This involved setting up of a downward flow in the water channel (see section 1.3) against which the buoyant particles rose. 3D trajectories of the particles were obtained using an orthogonal camera arrangement. We found that even a marginal reduction in the particle density from that of the fluid can lead to extremely high accelerations for finite-sized particles. As in Chapter 2, these modifications in acceleration statistics occurred due to the vertical motion of these particles through the turbulent flow. However, here, the intense accelerations had origin in the unsteady forcing due to the vortices shed in the particle wake. By exploring a range of Ξ for marginally buoyant particles ($\Gamma \approx 0.92$), we showed that at sufficiently small Ga (≤ 30), the particles followed the turbulent flow fluctuations. With increasing Ga ($225 - 900$), the contribution from vortex-shedding becomes increasingly dominant. The behavior in this regime could not be captured by point-particle simulations, even after the inclusion of finite-size correction terms.

In Chapters 2 and 3, the focus was limited to the 3D translational dynamics of buoyant particles. However, a complete description of particle dynamics requires, in addition, the three rotational components. Motivated by this, we have described a simple and robust experimental technique to track simultaneously the position and orientation of a finite-sized buoyant sphere in turbulence (see Chapter 4). The method employs a minimization algorithm to obtain the orientation from the two dimensional projection of a specific pattern drawn onto the surface of the sphere. The method returns an absolute orientation reference, and is easily scalable using parallel processing. Preliminary results validating the accuracy of the method using water tunnel experiments have also been presented.

In Chapter 5, we have tuned the rotational freedom of buoyant spheres by

changing their moment of inertia. We designed 3D-printed spheres of identical mass, but differing in their moments of inertia. Experiments in a turbulent flow have revealed that at reduced moment of inertia, rotation can couple with the particle's translational motion. This coupling can pave way for intense accelerations for a low moment of inertia spherical particle. Additionally, we observe that reducing moment of inertia at fixed mass density ratio can trigger a *tumble-flutter* transition for a buoyant sphere in turbulence. The tumbling spheres showed a tendency to drift towards the side-walls of the water channel, while the fluttering spheres remained in the bulk of the flow. We have extended the study to the case of freely rising buoyant spheres, i.e spheres rising through undisturbed fluid. Reducing the moment of inertia induces an increase in the oscillation amplitude, along with a change in the wake-pattern.

In Chapter 5, we had observed that changing the moment of inertia can trigger a change in the wake pattern of freely rising spheres. However, the experiments were limited to a few mass- and moment-of-inertia-ratios. Therefore, in Chapter 6, we have used direct numerical simulations to systematically explore the effect of mass and moment of inertia on the wake-induced dynamics of freely rising/falling circular cylinders. The amplitude of oscillation and the overall-system dynamics are linked to both the mass and the moment of inertia. We have revealed the existence of a mass dominated regime, which gradually transitions to a moment of inertia dominated regime. Large amplitude oscillations of rising cylinders fall within the moment of inertia dominated regime. The choice of body geometry (circular cylinder) and the state of the surrounding fluid (quiescent) are simplified. Nevertheless, this model problem has provided many insights into the underlying dynamics of buoyancy-driven bodies in general.

In Chapter 7, we studied the motion of heavy and buoyant cylindrical pendulums in quiescent water. This may be viewed as an extension of Chapter 6, but with a constraint on the cylinder. We experimentally measured the damped oscillations of the pendulum in a regime when the cylinder Reynolds number is large. The frequency of oscillations and the peak accelerations were shown to be dependent on the pendulum mass density ratio. We proposed a simple theoretical model, which takes buoyancy, fluid drag, added mass, and bearing friction into account. The model provides reasonable predictions for the frequency and peak accelerations. Using PIV, we measured the fluid flow surrounding the cylinder. The flow modifications by the cylinder during the first swing affects its subsequent oscillations. These aspects need to be included for improved predictions.

Chapters 2–7 addressed the dynamics of rigid buoyant particles. In Chapter 8, we move from rigid to deformable light particles, i.e a swarm of high

Reynolds number bubbles rising within an incident turbulent flow. The volume fractions reach up to 1%. At these concentrations, the bubble swarm may no longer be regarded as dilute. The liquid agitation in this case results from a complex interplay between the turbulence, the bubbles wakes, and the interaction between the bubble wakes. We have used the bubble parameter b [19, 112, 182], which measures the ratio of the bubble-induced kinetic energy to the kinetic energy of the turbulent liquid fluctuations before bubble injection, to characterize the bubbly flow. We have shown that the bubble parameter is a suitable parameter to capture the enhancement of liquid velocity fluctuations. Further, we have performed conditional measurements of the velocity downstream of the bubbles. This has allowed us to disentangle three regions of the wake, which have specific statistical properties. The primary wake region, close to the bubble, is similar to the wake of a single bubble rising in quiescent liquid. This is followed by a secondary wake, where the conditional probability density function of the velocity fluctuations is asymmetric and shows an exponential tail for the positive fluctuations, and a Gaussian one for the negative fluctuations. Further downstream, the liquid agitation is nearly Gaussian, and resembles a homogeneous and isotropic turbulent flow. The overall agitation in the bubbly turbulent flow is thus the result of these three contributions.

Future directions

This thesis has addressed a number of topics related to the dynamics of buoyant particles in turbulent flows. Many issues remain unexplored and open for investigation; some of these will be discussed below.

In Chapter 2, we restricted our analysis to particles with very small Stokes number. This means that they do not cluster in a turbulent flow. However, St is finite in many practical settings. One would then expect to observe particle clustering [32], wherein buoyant and heavy particles behave very differently due to the competing effects of buoyancy, added mass, and inertia. Thus, the effect of gravity on the behavior of finite St particles of arbitrary density will be an interesting issue for future study.

Chapters 2 & 3 dealt with two independent effects of gravity on buoyant particles in turbulence. It will be interesting to explore a regime where both the crossing-trajectory effect (Chapter 2) and the wake effect (Chapter 3) become important. We note that for a 1 mm bubble rising through the turbulent flow in the water channel, both effects will be important.

Chapter 5 introduced the effects of rotational inertia on the dynamics of buoyant spheres in turbulent and quiescent environments. The variation in parameter space was rather limited in this study. In future, we plan to explore a

wider parameter space for buoyant spheres rising in turbulence. In Chapter 5, we also found that a high moment of inertia sphere migrates towards the walls of the water channel flow. On similar lines, recent studies have shown that spherical bubbles migrate toward channel walls and could be used for drag reduction and heat transfer enhancement in vertical channel flows [203, 204]. However, this did not occur for deformable bubbles, and therefore, can be difficult to control in experiments. Our observation with spheres therefore opens up the possibility of designing particles which preferentially accumulate near the walls, with potential for drag and heat transfer modification in vertical channel flows.

In Chapter 6, we presented a numerical study of the motion and wakes of freely rising and falling cylinders at Galileo number $Ga = 500$. As a future step, it will be interesting to look at the effect of Ga on the same. Further, the Ga chosen in our simulations does not necessarily result in two-dimensional wakes. Therefore, fully three-dimensional simulations [77, 90, 205] would be useful to understand the possible effects of imposing two dimensionality in our simulations.

In Chapter 8, we confirmed that the -3 slope of energy spectrum of bubbly turbulence is robust to changes in turbulence and bubble volume fractions [35]. However, in most of the prior experiments, including ours, the bubble size lies in the range of $2 - 4$ mm. While a few studies looked into microbubbly turbulence at low volume fractions [192], the effect on the energy spectrum was not conclusive. An interesting question would be whether the -3 scaling is robust to changes in bubble size. Only after such an exploration, could one claim the -3 slope as a generic feature of bubbly turbulence.

Finally, we remind the reader that this thesis deals mainly with spherical and nearly spherical buoyant particles in turbulence. However, in nature and industry, one encounters anisotropic particles, which come in various shapes, and differ in density from the carrier fluid. Almost all the studies on anisotropic particles in turbulence have been numerical, and most have worked with the simplified case of inertial particles without gravity [206]. A few experiments addressed the motion of anisotropic particles in turbulence [207, 208], but these also ignored the effects of buoyancy. Once *buoyancy* is added to such a system, a variety of new physical phenomena can be expected. The Twente Water Tunnel facility is ideally suited to study this class of heavy/light anisotropic bodies that settle/rise through turbulent flows. With regard to tracking the particles, the method developed in Chapter 4 could be easily adapted to anisotropic shapes such as spheroids and ellipsoids. This can shed light on the dynamics of many naturally occurring objects, such as pollen, plant seeds, sediments, and plankton in the oceans.

Part I

Appendices

Gravity effects on small bubbles and particles in turbulence

In this appendix we provide some additional information about the numerical simulations and modelling from Chapter 2.

A.1 Acceleration flatness

Intermittency can be characterized in terms of the flatness of acceleration $\mathcal{F}(a_p) \equiv \langle a_p^4 \rangle / \langle a_p^2 \rangle^2$. Assuming statistical independence between $D_t u_i$ and $\partial_z u_i^*$, we obtain the tracer-normalized flatness of particle acceleration,

$$\frac{\mathcal{F}(\ddot{x})}{\mathcal{F}(D_t u_x)} = \frac{1 + \frac{12}{15 a_0 \mathcal{F}(D_t u_x)} \left(\frac{\text{St}}{\text{Fr}}\right)^2 \left(1 + \frac{\mathcal{F}(\partial_z u_x)}{45 a_0} \left(\frac{\text{St}}{\text{Fr}}\right)^2\right)}{1 + 2 \frac{2}{15 a_0} \left(\frac{\text{St}}{\text{Fr}}\right)^2 + \left(\frac{2}{15 a_0}\right)^2 \left(\frac{\text{St}}{\text{Fr}}\right)^4} \quad (\text{A.1})$$

$$\frac{\mathcal{F}(\ddot{z})}{\mathcal{F}(D_t u_x)} = \frac{1 + \frac{6}{15 a_0 \mathcal{F}(D_t u_x)} \left(\frac{\text{St}}{\text{Fr}}\right)^2 \left(1 + \frac{\mathcal{F}(\partial_z u_z)}{90 a_0} \left(\frac{\text{St}}{\text{Fr}}\right)^2\right)}{1 + 2 \frac{1}{15 a_0} \left(\frac{\text{St}}{\text{Fr}}\right)^2 + \left(\frac{1}{15 a_0}\right)^2 \left(\frac{\text{St}}{\text{Fr}}\right)^4} \quad (\text{A.2})$$

*The absence of correlations between $D_t u_i$ and $\partial_z u_i$ is a first order (to be refined) approximation.

In the limit of small St/Fr ,

$$\frac{\mathcal{F}(\ddot{x})}{\mathcal{F}(D_t u_x)} \simeq 1 - \frac{4}{15 a_0} \left(1 - \frac{3}{\mathcal{F}(D_t u_x)} \right) \left(\frac{\text{St}}{\text{Fr}} \right)^2 \quad (\text{A.3})$$

$$\frac{\mathcal{F}(\ddot{z})}{\mathcal{F}(D_t u_x)} \simeq 1 - \frac{2}{15 a_0} \left(1 - \frac{3}{\mathcal{F}(D_t u_x)} \right) \left(\frac{\text{St}}{\text{Fr}} \right)^2 \quad (\text{A.4})$$

It is verified that $\mathcal{F}(D_t u_x) > 3$. Therefore, both components are decreasing functions of St/Fr , with the vertical being larger than the horizontal one for small St/Fr .

In the large St/Fr limit,

$$\frac{\mathcal{F}(\ddot{x})}{\mathcal{F}(D_t u_x)} \simeq \frac{\mathcal{F}(\partial_z u_x)}{\mathcal{F}(D_t u_x)} \quad (\text{A.5})$$

$$\frac{\mathcal{F}(\ddot{z})}{\mathcal{F}(D_t u_x)} \simeq \frac{\mathcal{F}(\partial_z u_z)}{\mathcal{F}(D_t u_x)} \quad (\text{A.6})$$

It is verified that $\mathcal{F}(\partial_z u_z) < \mathcal{F}(\partial_z u_x) < \mathcal{F}(D_t u_x)$ [67]. This leads to the prediction $\mathcal{F}(a_v) < \mathcal{F}(a_h)$ i.e. vertical acceleration is less intermittent compared to the horizontal.

Figure A.1: Normalized flatness factor for small St/Fr . Hollow symbols show the flatness of horizontal acceleration. Solid symbols show the flatness of vertical acceleration. The vertical acceleration flatness is mostly higher in numerics, in qualitative agreement with the predictions of eq. (A.3)-(A.4).

Eq. (A.3)-(A.4) and eq. (A.5)-(A.6) provide predictions for the normalized acceleration flatness (a representation of intermittency) in the limits of small St/Fr and large St/Fr , respectively. Eq. (A.3)-(A.4) predict that the vertical component of flatness

exceeds the horizontal one in the small St/Fr limit, while eq. (A.5)-(A.6) predict that the horizontal component exceeds the vertical one when St/Fr is large. Therefore, an interesting cross-over is predicted between the flatness factors of the two components as one moves from small St/Fr to large St/Fr . In Fig. A.1, we present our numerical results in the small St/Fr range ($St/Fr < 1$). Despite the scatter in data, we make some interesting observations. With the exception of a single datapoint, the vertical components (solid symbols) are always higher as compared to the horizontal components (hollow symbols), in agreement with eq. (A.3)-(A.4). In the large St/Fr limit, as was clear from Fig. 2.5(c) of Chapter 2, the horizontal component exceeded the vertical one, again in agreement with the predictions (eq. (A.5)-(A.6)). Therefore, the cross-over predicted by us is qualitatively seen in the simulations as well. A quantitative agreement is missing since the higher moments (Flatness) are very sensitive. Moreover, the lowest Stokes numbers possible in simulations is ≈ 0.05 , which is still an approximation of the $St \rightarrow 0$ limit. In addition, the predictions are subject to a few assumptions, such as the absence of correlation between velocity field and its gradient. At present, we do not have the resolution to verify the intricate details of (A.3)-(A.4). These aspects could become the subject of future studies.

B

Finite-sized spheres in turbulence

In this appendix we discuss some of the experimental details of Chapter 3.

B.1 Particle detection and tracking

Figure B.1: (a) A sample original image with overlapping circles (all black). This is a rather rare situation where the spheres are very close to each other. (b) Circle detected image.

The circular Hough transform method is used to detect circle-centers and radii from the images (see Fig. B.1). Even cases with overlapping spheres are detected accurately. The sphere radius in the image varies a bit due to perspective effects, which is usually corrected in the calibration step by employing a circle-radius specific pixel to millimeter conversion. Two-dimensional trajectories are subsequently obtained by using a particle-tracking algorithm. These trajectories are subjected to smoothing [95], and 2D velocities and accelerations are obtained by taking the derivatives of the smoothed particle positions. Rare incidences of particle-particle collisions

are eliminated in the data processing stage by removing trajectories where the mean separation distance between spheres is within two diameters [81]. The vertical components of particle acceleration are correlated between the two cameras. A correlation coefficient ≥ 0.98 suggests that the trajectories belong to the same particle. Additional checks are performed by also correlating the vertical velocities, and by ensuring that the particle positions are matched in the two images. The third component extracted from Camera 2 is appended to the trajectory from Camera 1. The 3D trajectories of spheres with increasing Galileo numbers ($\text{Ga} \approx 30, 900, 2000$, and 3000) are provided as Supplementary Movies A, B, C, and D in [96–99]. The videos highlight the rich variability that can occur in buoyant-particle-dynamics in turbulence depending on the strength of the particle wake.

B.2 Lagrangian time correlation

Figure B.2: (a) Horizontal and (b) vertical components of Lagrangian temporal autocorrelation function of velocity and acceleration for the *very buoyant* ($\Gamma \approx 0.02$) sphere at $\text{Ga} \approx 3000$. The horizontal components have a period of oscillation double that of the vertical components. The period of oscillation is described well by the vortex shedding time scale τ_v . The acceleration autocorrelations are positively offset by 2 for clarity.

In Fig. B.2, we plot the Lagrangian time correlation function of velocity and acceleration for the ‘very buoyant’ ($\Gamma \approx 0.02$) spheres with Galileo number (Ga) ≈ 3000 . The left and right figures correspond to the horizontal and vertical components, respectively. The horizontal velocities and accelerations have the same oscillation frequency. The period of oscillation in the horizontal direction is double that in the vertical direction (right figure). This is typical of vortex-induced motions for rising and falling bodies in a quiescent fluid. In the supplementary Movie D [99], we show a typical trajectory traced by this particle at $Re_\lambda \approx 300$. The trajectory is color-coded by the instantaneous acceleration magnitude normalized by its standard deviation.

Controlling rotation

In this appendix we provide some details about the analysis presented in Chapter 5.

C.1 Tumbling vs fluttering dynamics

Fig. C.1(A) shows the Lagrangian autocorrelation function of velocity $C_v(\tau)$ for the low moment of inertia (MoI) sphere in turbulent flow. $C_v(\tau)$ shows strong oscillations for both horizontal (*red*) and vertical (*blue*) components of velocity. The frequency of horizontal fluctuations is half that of the vertical fluctuations and matches roughly the vortex shedding frequency $f_{viv} = 1/\tau_{viv}$. The same behavior is observed for the high MoI sphere, however with a marginal decrease in f_{viv} . Now we compare the Lagrangian autocorrelation of a horizontal component of angular velocity $C_\omega(\tau)$ (see Fig. C.1)(B). For the low MoI sphere (*red*), we see oscillations at a frequency which exactly matches the frequency of the translational velocity in Fig. C.1(A). Thus, the rotation appears to be synchronized with the translation of the sphere. However, for the high MoI sphere, the rotational autocorrelation $C_\omega(\tau)$ gradually decreases to zero in a timescale significantly larger than the vortex shedding timescale. The periodicity also disappears. Thus the translation and rotation do not appear to be synchronized for the high MoI sphere.

The mean horizontal drifting rates for the two spheres were shown in Fig. 5.3(A). Each datapoint corresponds to a separate experimental run. Between the two spheres, there is a significant difference in the horizontal drift speed. The high MoI sphere shows a horizontal drift that is nearly five times that of the low MoI sphere. From Movie S1 [145], we observe that the high MoI sphere tumbles in the flow with respect to this mean drifting direction. Statistical evidence for this may be seen when we compare the horizontal drifting direction with the direction of rotation. In Fig. 5.3(B), we plot the angle between $\vec{\omega} \times \hat{j}$ and the mean horizontal drift direction. Here, $\vec{\omega}$ is

Figure C.1: (A) Lagrangian temporal autocorrelation of the horizontal (red) and vertical (blue) velocity components of the low MoI sphere in turbulent flow. The horizontal axis is normalized by the vortex-shedding based time-scale $\tau_{viv} \propto D/U$, where D is the diameter of the sphere and U is the mean rise velocity of the particle. The vertical velocity fluctuations have twice the frequency of the horizontal velocity fluctuations, a well-known characteristic of vortex-induced vibrations [25]. (B) Lagrangian temporal autocorrelation of a horizontal component of angular velocity for the high (green) and low (red) MoI spheres in turbulent flow. For the low MoI sphere, the angular velocity fluctuations occur at the same frequency as the linear velocity fluctuations shown in (A), i.e the fluttering motion occurs in synchronization with the rotational motion. For the high MoI sphere, the angular velocity decorrelates slowly and without periodicity.

the angular velocity vector, and $\hat{j}$ is the unit vector pointing vertically upwards. 0° indicates alignment, and 180° indicates anti-alignment. The high MoI sphere shows alignment between the two vectors, indicating that the sphere (statistically) tumbles in the flow. Upon reducing the moment of inertia, the sphere flutters in the flow, as signified by the two peaks around 0° and 180° , and also evident from Movie S2 [146]. We believe that the tumbling motion of the high MoI sphere aids it to drift horizontally in the flow, while the fluttering motion of the low MoI sphere stabilizes it to remain in the middle of the measurement section. A direct analogy may be drawn to the motions reported for falling disks, strips and paper, where a moment-of-inertia reduction triggers a *tumble-flutter* transition [148]. A geometrical anisotropy of the rising/falling particle was considered necessary to induce this transition. The present observation for spheres suggests that this geometrical anisotropy factor may not be necessary for inducing this transition.

C.2 Dye visualization experiment

Fig. 5.5 showed the experimental setup used for visualizing the trajectories and wakes of the freely rising spheres. This is again shown below in Fig. C.2. The setup consists of a glass tank with dimensions $280 \text{ mm} \times 280 \text{ mm}$ wide, and 1500 mm high. The tank was filled with either water or a mixture of glycerine and water. The viscosity of the mixture was monitored at regular intervals using a rheometer. The sphere was held to the bottom of the tank by applying suction, and fluorescent dye (fluorescein-sodium) was injected into the cylindrical container just above the sphere. The dye was prepared

Figure C.2: Experimental arrangement used for visualizing the trajectories and wakes of freely rising spheres. Fluorescein sodium dye at 5×10^{-3} mol/l concentration is injected at the bottom of the water tank. The sphere is held at the bottom by applying suction.

with a concentration $\approx 5 \times 10^{-3}$ mol/L. The dye has a molecular diffusion coefficient in water, $D_m = 4 \times 10^{-10}$ m²/s at 22°C, which is much smaller than the water kinematic viscosity ν . This leads to a high Schmidt number, $Sc \equiv \nu/D_m \approx 2500$ [132]. The tank was left undisturbed for one hour. The glass tank was illuminated by eight blue LED floodlights, with a wavelength ≈ 470 nm. This corresponds to a peak in the emission spectrum of the fluorescent dye. A camera was positioned facing a side-wall of the glass tank. A light filter placed in front of the camera ensured that the incident blue light from the LED lamps was filtered out. The sphere was released by gently withdrawing the suction. The sphere rose through the liquid and entrained the dye on its way up. The entrained dye in the wake was slowly shed as the sphere rose, enabling us to visualize the wake. This gave a qualitative picture of the wake structure. The experiment was performed at two Galileo numbers: $Ga \approx 6000$ and $Ga \approx 500$. At high Ga , the wakes diffuse quickly, while at lower Ga , the vortex rings shed by the wakes are clearly visible. Movies S4 & S5 [151, 152] show the wake patterns at these Ga .

C.3 Movies

C.3.1 1. Movie S1 [145]

A typical trajectory of the high moment of inertia (MoI) sphere in a turbulent flow. The sphere tumbles in the flow, which is believed to aid the horizontal drift in the flow. Note that a mean downward flow roughly matches the sphere rise velocity enabling us to observe these motions for long duration.

C.3.2 2. Movie S2 [146]

A typical trajectory of the low moment of inertia (MoI) sphere in a turbulent flow. The sphere flutters in the flow, which stabilizes it to have a mean vertical rise trajectory. As a result of this, the particle remains in the bulk of the flow. Note that a mean downward flow roughly matches the sphere rise velocity enabling us to observe these motions for long duration.

C.3.3 4. Movie S4 [151]

Free-rise trajectories and wakes of the high MoI (*left*) and low MoI (*right*) spheres rising in still water at $Ga \approx 6000$.

C.3.4 5. Movie S5 [152]

Free-rise trajectories and wake of high MoI (*left*) and low MoI (*right*) spheres rising in 2.3:1 by volume mixture of glycerine and water at $Ga \approx 500$.

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Summary

The ubiquity of particle-laden turbulent flows in nature and technology makes it important to study them. This thesis covers mainly the broad topic of buoyant particles in turbulent flows. These particles rise through the carrier fluid, and their rising motion can lead to a variety of interesting physical phenomena. We have carried out experiments in a unique large-scale turbulence facility: the Twente Water Tunnel (TWT), which is part of the European High-performance Infrastructures in Turbulence (EuHIT) consortium.

We first addressed the acceleration statistics of small buoyant particles in turbulence (Chapter 2). To this end, we injected microbubbles (diameter $\sim 100 \mu\text{m}$) in the Twente Water Tunnel flow and tracked their motion using a moving camera setup. We compared the accelerations of the microbubbles to that of similar sized neutrally buoyant tracer particles. This provided the first evidence that tiny bubbles do not necessarily behave like fluid tracers, when it comes to the statistics of acceleration. Using numerics and theory, we showed that the deviations from neutrally buoyant tracers occur due to the finite drift velocity of these bubbles through the turbulent flow. We obtained quantitative predictions for the variance, time-correlation, and intermittency of acceleration. This has provided the first demonstration of the effects of crossing trajectories on the acceleration statistics of buoyant particles in turbulence.

As a next step, we studied the dynamics of finite-sized buoyant particles in a turbulent water flow (Chapter 3), i.e particles having a diameter larger than the smallest length scale of the flow, and with a density that is lower than that of water. We found that the larger the size of the particle, the more intense are the accelerations it experiences. This appears to be in contradiction to inertial range scaling predictions, which suggest that with increasing size, the acceleration variance ought to decrease as $d_p^{-2/3}$, where d_p is the particle diameter. We showed that the apparent reversal of trend occurs due to the wake-induced forces on the particles. With increasing particle size, while the forces due to turbulence decrease, the unsteady wake-induced forces on the particle increase. The latter outweighs the contribution from turbulence, resulting in an overall increase in the particle's acceleration variance.

In Chapter 3, we obtained the three dimensional trajectories of spheres moving in a turbulent flow. However, a full description of particle motion requires knowledge of their rotation as well. In Chapter 4, we described a method to obtain the three degrees of rotation of the spherical particle. The method employs a minimization algorithm to obtain the orientation from the 2D projection of an analytically prescribed pattern drawn onto the surface of the sphere. This method has been used to track the translation and rotation of buoyant spheres in turbulence in Chapter 5. Here, we showed that tuning the rotational inertia of a buoyant sphere can change its overall dynamics. A low moment of inertia sphere experiences intense accelerations when compared to a high moment of inertia sphere. We reveal that this occurs due to the increased rotational motion, which presumably changes the wake-induced forces on the sphere. This highlights the particle moment of inertia as an important control parameter governing the dynamics of buoyant spherical particles in flows. In Chapter 6, motivated by our observations in Chapter 5, we have conducted a numerical study on the effect of mass and moment of inertia on the two-dimensional motion and wakes of freely rising/falling circular cylinders. Here again, we found that both mass and moment of inertia control the dynamics of buoyant cylinders. Changing the moment of inertia affects the rotation of the particle. This has provided the first evidence that auto-rotation can significantly influence the trajectories and wakes of freely rising isotropic bodies.

In Chapter 7, we studied the dynamics of buoyant and heavy cylindrical pendulums undergoing large amplitude oscillations in water. By accounting for the fluid forces due to buoyancy, added mass, and (non-linear) drag, we are able to obtain a fair agreement with the experimental measurements of the pendulum motion.

Chapters 2–7 had addressed the dynamics of rigid buoyant particles. In Chapter 8, we moved from rigid to deformable light particles, i.e a swarm of 2 – 4 mm diameter air bubbles rising within an incident turbulent flow. The bubbles have a strong wake, and the liquid agitation induced by them is significant at the studied volume fractions ($\phi_V \sim 1\%$). We have used the *bubble* parameter b to characterize the turbulent bubbly flow. The *bubble* parameter b measures the ratio of the bubble-induced kinetic energy to the kinetic energy of the incident turbulent flow. Further, through conditional measurements of the velocity downstream of the bubbles, we have disentangled three regions of the flow: the primary wake, the secondary wake, and the far field. These have specific statistical properties. The overall agitation in the turbulent bubbly flow is the result of these three contributions.

Samenvatting

Turbulente stromingen met deeltjes zijn alomverteenwoordigd in natuur en technologie. Het is daarom van groot belang om de fysische aspecten ervan te bestuderen. Dit proefschrift omvat voornamelijk werk met betrekking tot lichte deeltjes in turbulente stromingen. Deze deeltjes stijgen op in de vloeistof en hun opwaartse beweging kan leiden tot een verscheidenheid aan interessante fysische verschijnselen. Wij hebben de experimenten uitgevoerd in een unieke grootschalige turbulentieopstelling: het Twente Waterkanaal (TWT), welk deel uitmaakt van het consortium Europese hoogwaardige infrastructuur voor turbulente stromingen (EuHIT).

We zijn begonnen met het onderzoeken van versnellingsstatistiek van kleine lichte deeltjes in turbulente stromingen (hoofdstuk 2). Hiervoor injecteerden we microbellen met een diameter van gemiddeld $100\ \mu\text{m}$ in het Twente Waterkanaal. De beweging van de microbellen werd vastgelegd met behulp van meerdere gemotoriseerde hogesnelheidscamera's. Hierdoor hebben wij de versnelling van de microbellen kunnen vergelijken met gelijkvormige deeltjes met neutraal drijfvermogen. Dit leidde tot het eerste bewijs dat kleine bellen zich niet noodzakelijkerwijs gedragen als deeltjes die de vloeistof perfect volgen (perfecte volgers) met betrekking tot de statistiek van de versnelling. Met behulp van computersimulaties en theorie hebben wij aangetoond dat een eindig snelheidsverschil van deze bellen ten opzichte van de turbulente stroming zorgt dat ze afwijken van perfecte volgers. We hebben kwantitatieve voorspellingen verkregen voor de variantie, tijdcorrelatie, en de intermittentie van de versnelling. Dit heeft voor de eerste keer laten zien dat kruisende banen van deeltjes invloed heeft op hun versnellingsstatistiek.

De volgende stap was het bestuderen van de dynamica van grote lichte deeltjes in een turbulente stroming (hoofdstuk 3). Dit zijn deeltjes met een diameter groter dan de kleinste schaal van de stroming en met een kleinere dichtheid dan water. Wij ontdekten dat hoe groter de diameter van het deeltje, des te intenser zijn de versnellingen die ze ervaren. Dit lijkt in tegenspraak te zijn met de schalingsvoorspelling in het inertiaële bereik, welk suggereert dat bij toenemende diameter de versnellingsvariantie moet dalen als $d_p^{2/3}$, waar d_p de diameter van het deeltje is. We hebben aangetoond dat de schijnbare omkering

van de trend zich voordoet door de door het zog veroorzaakte krachten op de deeltjes. Met toenemende deeltjesgrootte nemen de krachten als gevolg van turbulentie af, echter nemen de onstabiele zog-geïnduceerde krachten op het deeltje toe. Deze krachten zijn groter dan die van de turbulente stroming, wat resulteert in een algehele toename van de versnellingsvariantie van het deeltje.

In hoofdstuk 3 verkregen wij driedimensionale banen van de lichte deeltjes die zich in een turbulente stroming bewegen. Een volledige beschrijving van de deeltjesbeweging vereist echter ook kennis van hun rotatie. In hoofdstuk 4 beschrijven we een methode om de drie rotatiegraden van een bolvormig deeltje te verkrijgen. De methode maakt gebruik van een minimaliseringssgoritme om de oriëntatie te verkrijgen door middel van de een 2D-projectie van een analytisch voorgeschreven getekend patroon op het deeltje. Deze methode is gebruikt in hoofdstuk 5 om de translatie en rotatie te verkrijgen van de lichte deeltjes in de turbulente stroming. Hiermee hebben wij aangetoond dat het traagheidsmoment van een licht deeltje zijn dynamiek kan veranderen. Een deeltje met een laag traagheidsmoment ervaart intensere versnellingen vergeleken met deeltjes met een hoog traagheidsmoment. We laten zien dat dit gebeurt door de verhoogde rotatiebeweging, die vermoedelijk de zog-geïnduceerde krachten op het deeltje verandert. Dit onderstreept dat het traagheidsmoment van het deeltje een belangrijke parameter is in de dynamica van lichte bolvormige deeltjes in turbulente stromingen. In hoofdstuk 6, gemotiveerd door onze observaties in hoofdstuk 5, hebben we een numerieke studie uitgevoerd over het effect van massa en traagheidsmoment op de tweedimensionale bewegingen en de zog van stijgende en vallende cirkelcilinders. Ook hier ontdekten we dat zowel de massa als het traagheidsmoment invloed hebben op de dynamica van lichte deeltjes. Het variëren van het traagheidsmoment heeft invloed op de rotatie van het deeltje. Dit heeft het eerste bewijs geleverd dat de eigen rotatie aanzienlijk de baan en de zog van stijgende isotrope lichamen kan beïnvloeden.

In hoofdstuk 7 bestudeerden we de dynamica van lichte en zware cilindrische slingers die grote amplitudeoscillaties ondergaan in water. Door rekening te houden met de verschillende krachten o.a. drijfvermogen, toegevoegde massa en (niet-lineaire) weerstand, kunnen we een goede overeenkomst vinden met de experimentele metingen van de slingerbeweging.

In hoofdstukken 2–7 hebben de dynamica van starre lichte deeltje behandeld. In hoofdstuk 8 gaan we in op vervormbare lichte deeltjes (i.p.v. starre deeltjes), d.w.z. een zwerm van opstijgende luchtbellens met een diameter tussen 2–4 mm in een turbulente stroming. De bellens hebben een sterke zog, en de agitatie door de vloeistof die daardoor wordt geïnduceerd is significant voor de door ons bestudeerde volumefracties ($\phi_V \sim 1\%$). We hebben de bubble

parameter b gebruikt om de turbulente bellenstroming te karakteriseren. De bubble parameter b is de verhouding van de kinetische energie in de zog van de bel en de kinetische energie van de turbulente stroming. Verder hebben we door middel van voorwaardelijke snelheidsmetingen stroomafwaarts van de bellen, de stroming in drie soorten kunnen opdelen: de primaire zog, de secundaire zog en het verre veld. Deze hebben elk karakteristieke statistische eigenschappen. De totale agitatie in de turbulente bellenstroming is het gevolg van deze drie bijdragen.

Scientific output

Publications

1. **V. Mathai**, V. N. Prakash, J. Brons, C. Sun, D. Lohse
Wake-driven dynamics of finite-sized buoyant particles in turbulence,
Phys. Rev. Lett. 115, 124501 (2015),
2. **V. Mathai**, E. Calzavarini, J. Brons, C. Sun, D. Lohse
Microbubbles and microparticles are not faithful tracers of turbulent acceleration,
Phys. Rev. Lett. 117, 024501 (2016),
3. **V. Mathai***, X. Zhu*, C. Sun, D. Lohse
Mass and moment of inertia govern the transition in the dynamics of freely rising cylinders,
(Under review)
4. **V. Mathai***, X. Zhu*, C. Sun, D. Lohse
Controlling particle path-instabilities by tuning rotational inertia,
(Under review)
5. **V. Mathai**, R. N. Govardhan, V. H. Arakeri
On the impact of a concave-nosed axisymmetric body on a free surface,
Appl. Phys. Lett. 106, 064101 (2015),
6. **V. Mathai**, M. W. M. Neut, E. P. van der Poel, C. Sun
Translational and rotational dynamics of a large buoyant sphere in turbulence,
Exp. Fluids 57, 51, (2016),
7. E. O. Almeras, **V. Mathai**, D. Lohse, C. Sun
Experimental investigation of the turbulence induced by a bubble swarm rising within an incident turbulence,
(Under review)
8. E. G. Jebbink, **V. Mathai**, J. T. Boersen, C. Sun, C. H. Slump, P. C. J. M. Goverde, M. Versluis, M. M. P. J. Reijnen
Hemodynamic comparison of stent configurations used for aortoiliac occlusive disease,
J. Vasc. Surg. (2016) (In press),

Manuscripts under preparation

9. **V. Mathai***, L. Loeffen*, T. Chan, S. Wildeman
Dynamics of underwater pendulums
10. B. Gvozdic, E. O. Almeras, **V. Mathai**, D. P. M. van Gils, C. Sun, D. Lohse
Heat transport enhancement in bubbly flows
11. E. O. Almeras, **V. Mathai**, D. Lohse, C. Sun
Mixing induced by bubbles within an incident turbulent flow
12. D. Bakhuis, R. A. Verschoof, **V. Mathai**, C. Sun, D. Lohse
Drag modification by particles in turbulence

Invited and selected talks

- Hannover Messe - Wissenstag, Germany (*Invited*), (April 2015)
- COST Action, Flowing matter conference, Ankara, Turkey (*Invited*), (April 2015)
- 69th APS Division of Fluid Dynamics, Portland, Oregon, USA. (November 2016)
- 68th APS Division of Fluid Dynamics, Boston, USA. (November 2015)
- Physics at FOM, Veldhoven, (2015)
- European Turbulence Conference (ETC) at Delft, Netherlands, (August 2015)
- International Conference on Multiphase Flow (ICMF) at Florence, Netherlands, (May 2016)

Selected student supervision

- Master students:
 - J. Brons (Graduated 09/2014)
 - P. van der Plas (Expected 06/2017)
- Bachelor Students:
 - Fedde Gaastra & Jaap Nieland (Graduated 06/2014)
 - Laura Loeffen (Graduated 05/2016)
 - Timothy Chan, Internship student (06/2016 – 08/2016)

Teaching experience

- Teaching assistant for the Master course: "Experimental techniques in physics of fluids" for the study Applied Physics in Twente. (2014–2015)

- Teaching assistant for the Master course: "Experimental techniques in physics of fluids" for the study Applied Physics in Twente. (2015–2016)
- Lab lectures in various courses for CTA, LDA, PTV, and PIV techniques.

Peer review

- Reviewer for Journal of Fluid Mechanics (JFM), Physical Review Letters (PRL), International Journal of Multiphase Flow (IJMF), Journal of Turbulence (JoT), and Chemical Engineering Science (CES).

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