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Settling-driven Rayleigh–Taylor instability in a pre-existing turbulent field

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Vertical particle transport controls environmental processes ranging from sediment deposition and the export of organic material and microplastics to the fate of minerals released for carbon dioxide removal. When a particle-laden layer overlies denser water, particles destabilize the interface and descend collectively as convective plumes, far faster than individual settling. Yet this mechanism has commonly been studied in quiescent fluids, whereas the surface ocean is turbulent. We use direct numerical simulations to examine how pre-existing turbulence modifies each stage of settling-driven convection, varying the turbulence amplitude $\mathcal{A}$ and particle diameter D . Stronger pre-existing turbulence advances the onset of the Rayleigh–Taylor instability by increasing the effective seed amplitude. This earlier onset leaves less time for sediment to accumulate, resulting in a shallower nose and a weaker buoyancy gradient when growth begins. Consequently, the growth rate is reduced and the fastest-growing and marginal wavenumbers shift to lower values. Because the buoyancy gradient evolves almost identically at a given particle size, pre-existing turbulence controls when the instability begins rather than the interface it grows from. The disturbance also loses coherence, the vertical velocity field being rearranged into filaments within one e-folding at large $\mathcal{A}$. In the nonlinear phase, increasing $\mathcal{A}$ weakens both the sediment flux and the effective settling velocity, with stronger effects for smaller particles. For $D = 10 \mu\text{m}$, the settling velocity decreases by a factor of 2 from the weakest to the strongest $\mathcal{A}$, whereas for $D = 30 \mu\text{m}$, both bulk and frontal estimates are nearly $\mathcal{A}$ -independent.

Key words: Sediment-laden flow, Rayleigh–Taylor instability, Pre-existing turbulence

1. Introduction

The settling and vertical transport of particles in fluids are central to a broad range of environmental flows. Examples include the deposition of suspended sediment from river plumes in lakes, estuaries, coastal seas and glacial plumes (Bradley 1965; Hoyal *et al.* 1999;

Davarpanah Jazi & Wells 2016; Hill & Lintern 2021; Wells & Dorrell 2021; Sutherland *et al.* 2020), and the vertical transport of biogenic particles and microplastics in the ocean (Rahmani *et al.* 2022; Sutherland *et al.* 2023a,b; Semcesen *et al.* 2026). In such flows, particles need not settle independently at their quiescent Stokes velocity. As particles cross a density interface, their redistribution can create a gravitationally unstable region and drive convective sedimentation. The resulting convective transport can occur at rates substantially greater than the settling velocity of individual particles (Bradley 1965; Davarpanah Jazi & Wells 2016; Yang *et al.* 2025). Understanding the dynamics of this collective settling process is therefore important for predicting particle residence times, vertical particle fluxes and deposition in naturally stratified waters. The present study is motivated in part by a more recent application of this problem to ocean alkalinity enhancement (OAE), a marine carbon dioxide removal (mCDR) strategy. In OAE applications, finely ground alkaline minerals are introduced into the surface ocean and must remain suspended long enough to dissolve and increase seawater alkalinity (Köhler *et al.* 2013; Renforth & Henderson 2017; Eisaman *et al.* 2023; Zhou *et al.* 2025). Their residence time in the surface ocean may therefore depend strongly on settling-driven convection rather than on Stokes settling alone (Yang & Timmermans 2024).

Conventional estimates of residence time are based on the Stokes settling velocity of an isolated particle in quiescent fluid. This can overestimate the residence time. When particle-laden, buoyant fluid overlies denser saline water, gravitational settling transports particles across the density interface and creates a layer of excess density, a “nose region”, sandwiched between the salinity and sediment interfaces (Burns & Meiburg 2012, 2015; Davarpanah Jazi & Wells 2020). This layer is gravitationally unstable, and the resulting convective instability organizes the particle field into descending plumes whose vertical transport can exceed the Stokes settling velocity by one to two orders of magnitude (Hoyal *et al.* 1999; Blanchette & Bush 2005; Yu *et al.* 2014; Burns & Meiburg 2015; Lu *et al.* 2022). Two distinct instability mechanisms operate. When the sediment diffusivity is much smaller than that of salt and the two interfaces overlap, the system behaves as a double-diffusive fingering layer, producing narrow, vertically aligned fingers; the analogous dynamics in sheared double-diffusive layers have been examined by Smyth (2008); Smyth & Kimura (2011). When settling is fast enough that the unstable layer outgrows the salinity interface, the two interfaces decouple and the dominant mode becomes Rayleigh–Taylor (RT) like, characterized by overturning motions and broad, coherent plumes. Burns & Meiburg (2012) showed via linear stability analysis, and Burns & Meiburg (2015) confirmed with nonlinear direct numerical simulation (DNS), that the transition between these regimes is governed by the ratio $R_t = H/l_s$ of the nose-region height to the salinity interface thickness, with fingering dominating for $R_t \lesssim O(0.1)$ and RT-like behaviour for larger values.

Because R_t increases with the Stokes settling velocity, particle size plays an important role in determining the dominant instability. Scale analysis by Green (1987) suggested that double-diffusive effects may become important for particle diameters below approximately $2\ \mu\text{m}$, while three-dimensional simulations by Yu *et al.* (2014) showed that particles with $D \sim 2\text{--}4\ \mu\text{m}$ produce narrow “leaking” fingers, characteristic of double-diffusive convection. For larger particles, settling more rapidly establishes an unstable bulk density profile, giving rise to RT plumes. Consistent with Yu *et al.* (2014), the sharp-interface simulations of Shao *et al.* (2017) showed a similar particle-size-dependent regime transition, with RT convection becoming dominant for particles of order $10\ \mu\text{m}$ and larger. We focus on this latter, RT-dominated regime and reach it by construction rather than through particle size alone: the sediment and salinity diffusivities are set equal, and, as Yu *et al.* (2014) note, double-diffusive effects become weak when the two diffusivities are comparable, leaving settling-driven accumulation as the route to instability. In this regime, collective convective

motions can substantially enhance vertical particle transport relative to individual Stokes settling. This regime also overlaps with the particle-size range commonly considered for particle-based OAE (Köhler *et al.* 2013; Eisaman *et al.* 2023; Yang & Timmermans 2024; Yang *et al.* 2025).

A common feature of the DNS studies above is that the instability is seeded from an essentially quiescent base state, perturbed either by low-amplitude white noise or by the eigenfunction of the fastest-growing linear mode (Burns & Meiburg 2015; Yu *et al.* 2014; Yang *et al.* 2025). This isolates the instability and permits clean comparison with linear theory (Smyth & Carpenter 2019; Liu *et al.* 2024). Moreover, Dalziel *et al.* (1999) demonstrated that, for classical RT mixing, the late-time growth of the mixing zone is largely insensitive to the amplitude of such perturbations, and retains little memory of the initial condition. The ocean surface mixed layer, however, is not quiescent: it is continuously stirred by wind, waves, and convection, so that a settling-driven instability at its base develops within a pre-existing turbulent field rather than growing from infinitesimal noise. The distinction is not new in the context of stratified shear instability, where Kaminski & Smyth (2019) showed that pre-existing turbulence substantially modifies the growth, structure, and mixing efficiency of Kelvin–Helmholtz billows, such that the resulting instability differs from the one a quiescent base state would predict. Liu *et al.* (2022) further showed that the transition to turbulence in a stratified shear layer depends sensitively on the perturbation from which it develops, with small changes in the initial field producing substantially different turbulent outcomes. Characterizing that sensitivity is not our main goal here, but it informs the design of our parameter sweep. We vary $\mathcal{A}$ over several orders of magnitude so that the systematic dependence on amplitude is large compared with the scatter among individual realizations of the initial field. We verify this by repeating one case with a different set of random phases. It remains unknown how a background turbulent field alters the onset, structure, or sediment transport of RT plumes.

We initialize three-dimensional DNS of the two-layer particle-laden configuration with a fully developed turbulent velocity field of prescribed amplitude $\mathcal{A}$, in place of the infinitesimal perturbations used previously. Each simulation is initialized with a synthetic turbulent velocity field constructed with a Kolmogorov $k^{-5/3}$ amplitude spectrum, scaled by an amplitude $\mathcal{A}$ that sets the total kinetic energy of the field. The amplitudes considered correspond to turbulent kinetic energy dissipation rates spanning $\varepsilon \sim O(10^{-8})$ – $O(10^{-4})$ W Kg⁻¹, covering the range reported for the upper ocean and surface mixed layer (Thorpe *et al.* 2003; Burchard *et al.* 2008; Bluteau *et al.* 2011; Kaminski *et al.* 2021). The Fourier phases are assigned randomly, so the initial field reproduces the energy distribution of ocean turbulence but not the coherent structure of a fully developed turbulent flow. The amplitudes are therefore calibrated to observed mixed-layer energy levels rather than chosen arbitrarily. We additionally vary the particle diameter over a range relevant to OAE and natural estuaries, allowing us to determine how pre-existing turbulence modifies RT-driven settling and how this sensitivity depends on particle size, including whether particle sizes are affected differently by background turbulence.

Our objectives are twofold. First, we seek a mechanistic understanding of how and why pre-existing turbulence modifies settling-driven convection throughout its evolution, from linear growth on an evolving base state, through the loss of modal coherence, to the nonlinear plumes that transport sediment downward. Second, we ask whether the influence of pre-existing turbulence is equally important across particle sizes.

The remainder of the paper is organized as follows. §2 presents the governing equations, dimensionless parameters, and construction of the pre-existing turbulent field. §3 examines the linear phase, applying stability analysis to the evolving mean buoyancy profile to determine how pre-existing turbulence sets the onset time and, consequently, the growth

rate and the selected wavenumbers. §4 examines the extent to which the normal-mode description remains valid by quantifying the rate at which the disturbance loses its spatial coherence. §5 describes the three-dimensional evolution from the onset of the instability until sediment first reaches the bottom boundary. §6 quantifies the resulting transport in terms of the vertical sediment flux, turbulent diffusivities, effective settling velocity, and delivered mass. Finally, §7 summarizes the main findings and discusses the limitations of the present study.

2. Methodology

2.1. Mathematical model

We consider an initially two-layer system in which the buoyancy b depends on both salinity S and the sediment concentration c ,

$$b(S, c) = -g (1 + \beta \Delta S + \gamma \Delta c), \quad (2.1)$$

where β and γ are the haline and sediment contraction coefficients, respectively.

The evolution of the velocity field $\mathbf{u} = (u, v, w)$ is described by the incompressible Navier–Stokes equations under the Boussinesq approximation,

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

$$\frac{\partial \mathbf{u}}{\partial t} + (\mathbf{u} \cdot \nabla) \mathbf{u} = -\frac{1}{\rho_0} \nabla p + b(S, c) \hat{\mathbf{k}} + \nu \nabla^2 \mathbf{u}, \quad (2.3)$$

where ρ_0 is a reference density and g is gravitational acceleration, $\hat{\mathbf{k}}$ is the vertical unit vector, p is pressure and ν is the kinematic viscosity. The two scalar fields are governed by

$$\frac{\partial S}{\partial t} + (\mathbf{u} \cdot \nabla) S = \kappa_S \nabla^2 S, \quad (2.4)$$

$$\frac{\partial c}{\partial t} + (\mathbf{u} \cdot \nabla) c = \kappa_c \nabla^2 c + W_{st} \frac{\partial c}{\partial z}, \quad (2.5)$$

where κ_S and κ_c are the saline and tracer diffusivities and $W_{st} > 0$ is the Stokes settling velocity of an individual particle. Particles are assumed monodisperse and dilute, so that particle–particle interactions (flocculation, hindered settling) are neglected (Yu *et al.* 2014; Burns & Meiburg 2015; Yang *et al.* 2025). The system is characterized by four dimensionless groups (Burns & Meiburg 2015; Yang *et al.* 2025): the density ratio (R_ρ), the nondimensional settling velocity (W_p), the Schmidt number (Sc), and the diffusivity ratio (τ),

$$R_\rho \equiv \frac{\gamma \Delta c}{\beta \Delta S} = 0.5, \quad W_p \equiv \frac{W_{st}}{(\nu g'_c)^{1/3}}, \quad Sc \equiv \frac{\nu}{\kappa_S} = 9, \quad \tau \equiv \frac{\kappa_S}{\kappa_c} = 1, \quad (2.6)$$

where $g'_c = \gamma \Delta c g$ is the reduced gravity associated with the tracer loading alone. Only the tracer loading enters g'_c because the Rayleigh–Taylor instability (RTI) is driven by the density anomaly in the unstable nose region that develops beneath the interface. τ is set to unity, so that the tracer and salinity diffuse at the same rate, i.e. $\kappa_c = \kappa_S = \nu/Sc$. Only W_p varies between cases. Its denominator $(\nu g'_c)^{1/3}$ is the velocity at which buoyancy generated by the tracer loading balances viscous damping, so W_p measures the settling of an individual particle against the collective motion that loading can drive. Viscosity enters because it suppresses the smallest scales, allowing the nose to develop to a finite size before

Parameter Symbol	g	ν	γ	Δc	β	ΔS	h_b	L_x	L_y	L_z
Value	9.81	1.0×10^{-6}	2.2	0.0045	0.76	0.025	0.0165	0.223	0.149	0.297
Dimensions	m s^{-2}	$\text{m}^2 \text{s}^{-1}$	—	—	—	—	m	m	m	m

Table 1. List of constants used in this study.

becoming unstable rather than overturning as soon as a density inversion appears. The three particle sizes examined, $D = 10, 15$ and $30 \mu\text{m}$, correspond to $W_p = 0.026, 0.059$ and 0.235 . All are small, so convection rather than individual settling dominates the descent in every case, but the margin narrows by an order of magnitude across the range. The strength of the classical RTI is set by the Atwood number, $At \equiv \Delta b_c / (2g + 2\Delta b_S + \Delta b_c) = 0.0048$, where $\Delta b_c \equiv g\gamma\Delta c = g'_c$ and $\Delta b_S \equiv g\beta\Delta S$ are the buoyancy differences associated with the sediment loading and with the salinity contrast, respectively. At is held fixed across all cases. The physical constants held fixed across all simulations are listed in table 1. The physical parameters of the particles correspond to those of silicon carbide particles commonly used in laboratory experiments and olivine particles used in OAE studies (e.g. Davarpanah Jazi & Wells 2016; Sutherland *et al.* 2018; Yang & Timmermans 2024; Yang *et al.* 2025). Unlike Konopliv *et al.* (2018); Yang *et al.* (2025), we impose no mean velocity shear. In their place we introduce two additional parameters, the pre-existing turbulence amplitude $\mathcal{A}$, defined and constructed in §2.3, and particle diameter D .

2.2. Linear stability analysis

To evaluate the linear instabilities, the governing equations are linearized about the horizontally averaged buoyancy profile $\langle b \rangle_{xy}(z, t)$ obtained from the simulations, treated as a quasi-steady base state at each time t . These equations are then subjected to perturbations induced by small-amplitude, normal-mode disturbances proportional to the real part of $a(z) \exp(\sigma t + ikx)$, where $a(z)$ denotes the vertically varying, complex amplitude of any perturbation quantity, σ stands for the complex exponential growth rate, and k is the wavenumber in the streamwise direction. The normal-mode equations are discretized using a Fourier-Galerkin method, yielding a generalized matrix eigenvalue problem that is solved using standard methods (Smyth & Carpenter 2019; Lian *et al.* 2020).

2.3. Numerical model

We perform three-dimensional direct numerical simulations using `Oceananigans.jl` (Wagner *et al.* 2025), a GPU-accelerated finite-volume solver for the incompressible Boussinesq equations on a staggered Arakawa C-grid. Advection of both momentum and tracers is discretized with a centred fourth-order scheme. Time-stepping uses a third-order Runge–Kutta scheme; the pressure is obtained by a fast-Fourier-transform-based projection consistent with the periodic horizontal boundaries.

The computational domain is $\{L_x \times L_y \times L_z\} = \{0.223 \times 0.149 \times 0.297\}$ m, discretized on a uniform grid of $\{N_x, N_y, N_z\} = \{768, 512, 1024\}$. The resolution is constrained by the Kolmogorov scale $L_K = (\nu^3/\varepsilon)^{1/4}$ or, for $Sc > 1$, by the smaller Batchelor scale $L_B = L_K Sc^{-1/2}$. In the present study, our computational resolution is $\sim L_B$. This grid spacing and domain geometry are adopted from Yang *et al.* (2025), whose convergence tests apply directly to the present configuration. They established that halving the grid spacing leaves the volume-averaged settling speed essentially unchanged, and that extending the depth of the lower layer likewise leaves the results unaffected, so neither the resolution nor the proximity of the bottom boundary influences the transport diagnostics.

208 Periodic boundary conditions are imposed in both horizontal directions. The upper and
209 lower boundaries are free-slip with no normal flow,

$$\frac{\partial u}{\partial z} = \frac{\partial v}{\partial z} = 0, \quad w = 0 \quad \text{at } z = -L_z/2, L_z/2. \quad (2.7)$$

210 Salinity satisfies a no-flux condition at both boundaries, $\partial S/\partial z = 0$. For the tracer, the
211 total (settling plus diffusive) flux vanishes at the upper boundary while the diffusive flux
212 vanishes at the lower boundary (Yu *et al.* 2014; Yang *et al.* 2025),

$$W_{st}c - \kappa_c \frac{\partial c}{\partial z} = 0 \quad \text{at } z = L_z/2, \quad \frac{\partial c}{\partial z} = 0 \quad \text{at } z = -L_z/2. \quad (2.8)$$

213 The initial mean state consists of a sediment-laden upper layer above a saltier lower
214 layer, with both interfaces of half-thickness h_b centred at $z = z_0 = L_z/4$ (dashed curves in
215 figure 1),

$$\bar{S}(z, t = 0) = \frac{\Delta S}{2} \left[1 - \tanh\left(\frac{z - z_0}{h_b}\right) \right], \quad (2.9)$$

$$\bar{c}(z, t = 0) = \frac{\Delta c}{2} \left[1 + \tanh\left(\frac{z - z_0}{h_b}\right) \right], \quad (2.10)$$

$$\bar{\mathbf{u}}(z, t = 0) = \mathbf{0} + \mathbf{u}'. \quad (2.11)$$

219 This state is statically stable at $t = 0$; the RTI arises only after differential settling has built
220 an unstable nose region beneath the salinity interface (solid curve in figure 1b). Superposed
221 on equation (2.11) is either white noise or pre-existing turbulent velocity perturbation.

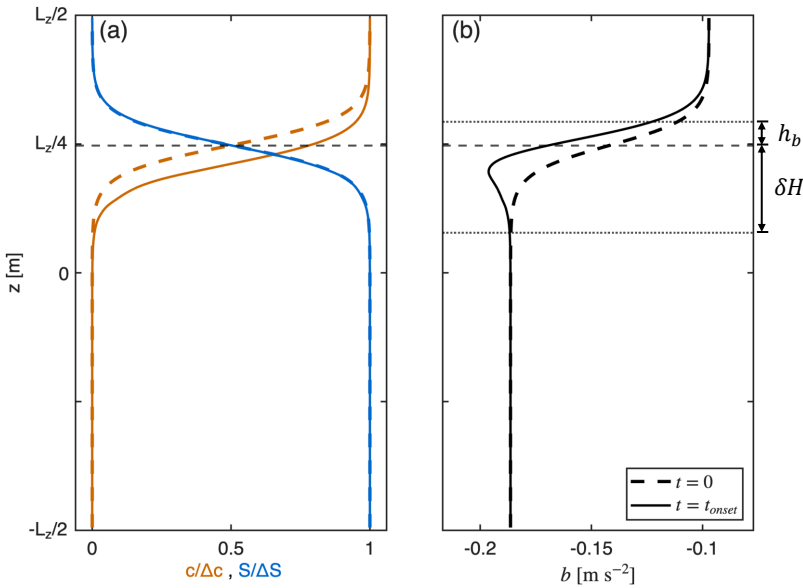


Figure 1. Vertical profiles at $t = 0$ (dashed) and at RTI onset t_{onset} (solid) of (a) normalized tracer concentration (brown) and salinity (blue), and (b) buoyancy b . The salinity profile remains nearly unchanged between the two times, while the tracer concentration profile descends, extending the buoyancy transition zone below the initial interface (dashed horizontal line at $L_z/4$). h_b denotes the half-thickness of the initial buoyancy interface, and δH is the additional nose thickness formed by the descending tracer by t_{onset} .

222 The pre-existing turbulent initial perturbation field is constructed in spectral space.
 223 We write the perturbation velocity as a Fourier series over the discrete wavenumbers
 224 $\mathbf{k} = (k_x, k_y, k_z)$ admitted by the domain,

$$u'_j(\mathbf{x}) = \text{Re} \left\{ \sum_{\mathbf{k}} \hat{u}_j(\mathbf{k}) e^{i\mathbf{k} \cdot \mathbf{x}} \right\}, \quad (2.12)$$

225 and construct the coefficients $\hat{u}_j$ directly. Let $k_h = (k_x^2 + k_y^2)^{1/2}$ denote the horizontal
 226 wavenumber magnitude. We first generate a provisional set of coefficients $\hat{u}_j^{(0)}$, in which
 227 each of the three velocity components is assigned a Kolmogorov-like amplitude in the
 228 horizontal and an independent random phase,

$$\hat{u}_j^{(0)}(\mathbf{k}) = \mathcal{A} k_h^{-5/6} e^{i\phi_j(\mathbf{k})}, \quad j = 1, 2, 3, \quad (2.13)$$

229 where the phases ϕ_j are drawn randomly in $[0, 2\pi)$. The spectral energy density, therefore,
 230 scales as $E \propto k_h^{-5/3}$. The spectral slope is imposed on the horizontal wavenumber only;
 231 this reflects the anisotropy of mixed-layer turbulence, in which the vertical scale is set by
 232 the layer depth while the horizontal scales follow an inertial-range cascade (Kaminski &
 233 Smyth 2019).

234 Incompressibility is enforced by projection onto the solenoidal subspace,

$$\hat{u}_j(\mathbf{k}) = \left(\delta_{jl} - \frac{k_j k_l}{|\mathbf{k}|^2} \right) \hat{u}_l^{(0)}(\mathbf{k}), \quad (2.14)$$

235 after which the physical-space field follows from equation (2.12). Note that the projection
 236 uses the full wavenumber vector, so the field remains three-dimensional and divergence-free
 237 despite the purely horizontal spectral weighting.

238 The relevance of the imposed turbulence to real ocean mixed-layer conditions is assessed
 239 through the volume-averaged dissipation rate prior to nose formation, which follows $\langle \bar{\varepsilon} \rangle_V =$
 240 $C_0 \mathcal{A}^2$ across the sweep, where $C_0 = 0.3384 \text{ s}^{-1}$ (figure 2) and spans 1.4×10^{-8} to
 241 $10^{-4} \text{ W kg}^{-1}$ for $\mathcal{A} = 2.07 \times 10^{-4}$ to 0.0172 m s^{-1} . This range sits within observed surface-
 242 ocean values (Zeiden *et al.* 2023; Esters *et al.* 2018; Iyer *et al.* 2021; Cen *et al.* 2022). Our
 243 lower amplitude cases ($\mathcal{A} \leq 4.15 \times 10^{-4}$; $\varepsilon < 10^{-7} \text{ W kg}^{-1}$) correspond to a weakly forced
 244 mixed layer, and our intermediate cases, where dissipation rate is $10^{-7} - 10^{-6} \text{ W kg}^{-1}$,
 245 correspond to moderate wind forcing. Strong amplitude case ($\mathcal{A} = 0.0172$) corresponds
 246 to the dissipation rate beneath actively breaking waves in a near-surface layer (Terray *et al.*
 247 1996), reaching $\varepsilon \sim 10^{-4} \text{ W kg}^{-1}$.

248 The full set of simulations is listed in table 2. The amplitudes are not sampled uniformly
 249 across the three particle sizes. Every particle-size family is run at $\mathcal{A} = 2.07 \times 10^{-4}$,
 250 4.15×10^{-3} and 1.72×10^{-2} , which spans nearly two decades and fixes the trend with
 251 amplitude. The $D = 15$ and $D = 30 \mu\text{m}$ families add further amplitudes within that interval.
 252 The $D = 15 \mu\text{m}$ case is thus the reference family, sampled finely enough to resolve the
 253 shape of the response, while the other two establish whether that response carries across
 254 particle size, for which the three shared amplitudes are sufficient. The response within the
 255 finely sampled interval is not monotonic, and appendix A examines what that implies. The
 256 last two columns, the onset time t_{onset} and the time at which sediment first reaches the
 257 bottom t_f , are outcomes rather than inputs; they are defined and analyzed in §3 and §5
 258 respectively, and are collected here for reference.

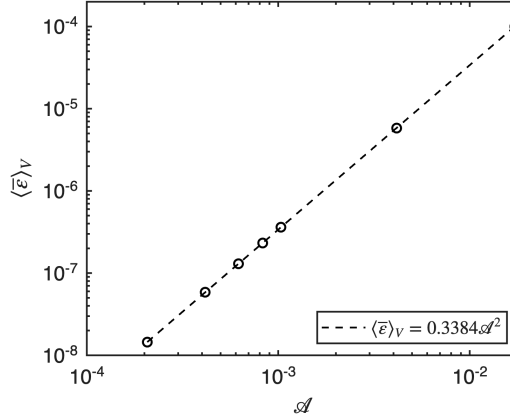


Figure 2. Volume averaged turbulent kinetic energy dissipation rate with time average before the nose region appears (from $t = 0$ to 21 s) with $D = 15 \mu\text{m}$.

Case	$D [\mu\text{m}]$	$W_{st} [\text{m s}^{-1}]$	$\mathcal{A} [\text{m s}^{-1}]$	Type	$t_{\text{onset}} [\text{s}]$	$t_f [\text{s}]$
1	10	1.2×10^{-4}	0.000207	P	79.4	102.8
2			0.00415	P	63.6	101.4
3			0.0172	P	53.8	91.8
4	15	2.7×10^{-4}	0.000207	P	40.0	59.6
5			0.000415	P	37.8	61.0
6			0.000621	P	36.8	61.4
7			0.00083	P	35.4	54.8
8			0.00135	P	35.2	59.8
9			0.00415	P	29.2	58.0
10			0.0172	P	22.8	54.6
11			0.0007	W	43.2	61.6
12			0.000415	W	42.2	63
13	30	10.8×10^{-4}	0.000207	P	15	28.6
14			0.000415	P	13.8	28.4
15			0.00083	P	12.8	26.8
16			0.00415	P	9.8	25.8
17			0.0172	P	7.4	25.8

Table 2. Simulation parameters and timing for all cases. D is the particle diameter and W_{st} the corresponding Stokes settling velocity; $\mathcal{A}$ is the amplitude of the initial velocity perturbation. “Type” distinguishes cases seeded with pre-existing turbulence (P) from those seeded with white noise (W). t_{onset} is the RTI onset time, defined as the first time the volume-averaged $|w'c'|$ exceeds 1% of its maximum, and t_f is the time at which sediment first reaches the bottom boundary. Cases are grouped by particle size. The two white-noise cases at $D = 15 \mu\text{m}$ are both compared against case 5: case 12 matches it in nominal amplitude $\mathcal{A}$, and case 11 in total kinetic energy, the two constructions differing in energy at equal $\mathcal{A}$ because of the spectral weighting in equation (2.13). Case 7 includes 3 ensembles.

3. Linear phase

We examine how pre-existing turbulence affects both the onset of RTI and its subsequent growth rate. We apply linear stability analysis to the horizontally averaged buoyancy profile extracted from the DNS at each output time step, treating the profile at each time step as a quasi-steady base state and predicting the instantaneous growth rate and corresponding

fastest-growing wavenumber, following approaches used in stratified shear instabilities (Carpenter *et al.* 2010; Smith *et al.* 2021; Liu *et al.* 2022; Smyth & Carpenter 2019). This quasi-steady treatment remains a reasonable approximation up to onset, (shown as dots in figure 3), beyond which nonlinear plume dynamics take over and the base state no longer evolves slowly relative to the perturbation growth. Unstable modes first appear once the

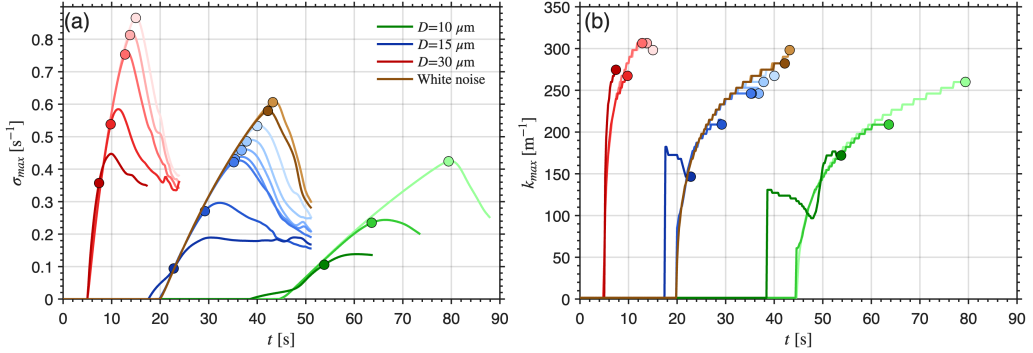


Figure 3. Time evolution of (a) the growth rate and (b) the fastest-growing wavenumber, obtained from linear stability analysis applied at each time step, for particle diameters $D=10, 15$ and $30 \mu\text{m}$ and for the white-noise initial condition (brown). Darker shades indicate larger $\mathcal{A}$, and filled circles mark the onset of RTI.

nose region emerges, where $\partial b / \partial z < 0$. Within each particle-size family, the time at which unstable modes first appear ($\sigma > 0$) varies little across turbulence amplitudes (figure 3a). We define the onset time t_{onset} as the first time at which the volume-averaged magnitude of the sediment flux, $\langle |w'c'| \rangle_V(t)$, exceeds 1% of its maximum value over the simulation (Yu *et al.* 2014; Yang *et al.* 2025), where prime is the deviation from the horizontal average. By this measure, t_{onset} is shorter for stronger pre-existing turbulence (figure 4; values for every case are given in table 2), and the effect of pre-existing turbulence on RTI onset is more significant for smaller particle size.

However, an earlier onset does not imply faster growth. Figure 3a shows that the fastest growth rate $\sigma_{\max}$ is systematically smaller for stronger pre-existing turbulence (darker colours denote larger $\mathcal{A}$), and the same ordering holds across particle sizes. The fastest-growing wavenumber $k_{\max}$ decreases in parallel (figure 3b).

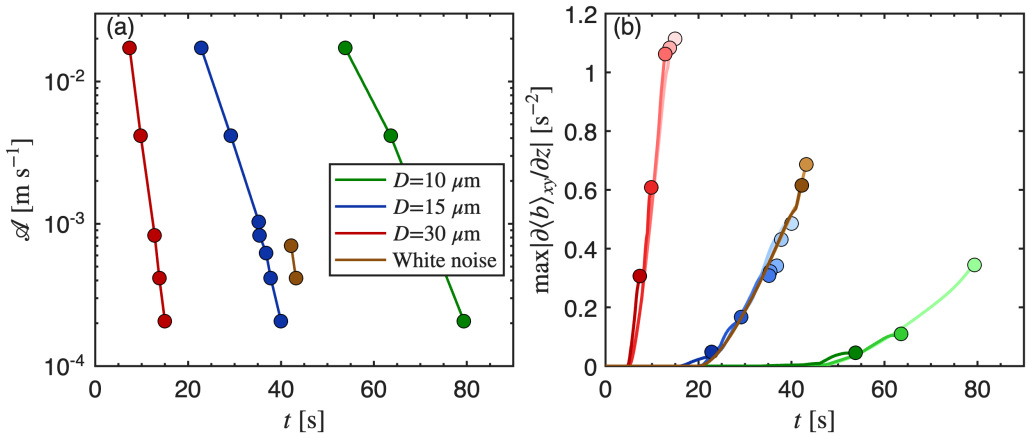


Figure 4. (a) Pre-existing turbulence amplitude $\mathcal{A}$ against the RTI onset time, coloured by particle diameter D . (b) Evolution of the peak buoyancy gradient for all cases, with filled circles marking each case's onset.

This suppression is consistent with the role of viscous and diffusive damping in setting the most unstable scale. In the inviscid limit, RTI growth rate increases monotonically toward the smallest scales; viscosity and diffusivity damp this tendency and select a finite fastest-growing wavenumber $k_{\max}$. Both trends originate in the buoyancy gradient of the base state at the time of onset. Figure 4(b) shows the evolution of $\max|\partial\langle b\rangle_{xy}/\partial z|$ for every case. Within each particle-size family the curves collapse onto a common trajectory, where the interface sharpens at a rate set by the Stokes settling velocity, independently of the pre-existing turbulence amplitude. Large particles settle faster and sharpen the interface sooner, which separates the families, but within a family, $\mathcal{A}$ does not alter how the gradient develops. What differs between cases is instead where along that trajectory onset occurs. Stronger pre-existing turbulence provides a larger initial disturbance, so the settling flux reaches the onset threshold sooner, at a stage when less sediment has accumulated beneath the interface and the nose is correspondingly less developed. The destabilizing buoyancy anomaly is therefore weaker at t_{onset} . Since the growth rate on a uniformly stratified layer scales as $\sqrt{-\partial b/\partial z}$ for $\partial b/\partial z < 0$, a weaker gradient limits $\sigma_{\max}$.

To provide a quantitative reference against which the diffusive suppression of small-scale modes can be assessed, we compare the numerically computed neutral wavenumber k^* , the value of k at which $\sigma(k, t_{\text{onset}})$ first crosses zero, against an analytical stability boundary derived from the viscous, diffusive dispersion relation for a normal-mode perturbation on a locally uniform buoyancy gradient (Smyth & Carpenter 2019, §2.3). For a perturbation $a(z) \exp(\sigma t + ikx)$ superposed on a uniformly stratified base state with gradient $\partial b/\partial z$, the growth rate satisfies

$$\sigma^2 + (\nu + \kappa_b)K^2\sigma + \nu\kappa_b K^4 + \frac{\partial b}{\partial z} \cos^2 \theta = 0, \quad (3.1)$$

where κ_b is the buoyancy diffusivity (since $\tau = 1$, $\kappa_b = \kappa_S = \kappa_c = \nu/\text{Sc}$). K is the wavenumber magnitude, $K = \sqrt{k^2 + \ell^2 + m^2}$ and θ is its inclination from the horizontal. The solutions for the growth rate are

$$\sigma = \frac{1}{2} \left[-(\nu + \kappa_b)K^2 \pm \sqrt{\mathcal{D}} \right], \quad (3.2)$$

where the discriminant is

$$\mathcal{D} = (\nu + \kappa_b)^2 K^4 - 4(\nu\kappa_b K^4 + \partial b/\partial z \cos^2 \theta). \quad (3.3)$$

The unstable case occurs when σ is real and positive, thus, $\mathcal{D} > (\nu + \kappa_b)^2 K^4$, i.e. $\nu\kappa_b K^4 < -\partial b/\partial z \cos^2 \theta$, which can be satisfied only when $\partial b/\partial z < 0$. Solving for K gives,

$$K < \left(\frac{-\partial b/\partial z \cos^2 \theta}{\nu\kappa_b} \right)^{1/4}. \quad (3.4)$$

The stability boundary ($\sigma = 0$) is evaluated at the orientation that maximizes growth: a purely horizontal wavevector ($\theta = 0$), which forces the associated fluid motion to be purely vertical. At $\theta = 0$ the vertical wavenumber vanishes, so K reduces to the horizontal wavenumber, we denote this stability boundary value

$$k_c = \left(\frac{-\partial b/\partial z}{\nu\kappa_b} \right)^{1/4}. \quad (3.5)$$

We evaluate k_c using the peak gradient, $\max(|\partial\langle b\rangle_{xy}/\partial z|)$ at t_{onset} . Although k_c is derived for a uniform buoyancy gradient, the peak value is the appropriate choice in this limit

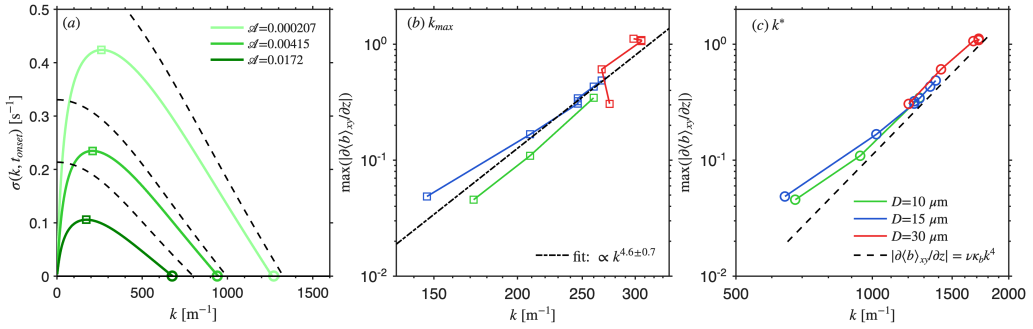


Figure 5. Linear stability of the horizontally averaged buoyancy profile at the onset of RTI. (a) Growth rate for the three pre-existing-turbulence amplitudes at $D = 10 \mu\text{m}$, with darker shading indicating larger $\mathcal{A}$. Squares mark the fastest-growing wavenumber k_{max} ; circles mark the neutral wavenumber k^* , at which σ crosses zero. Dashed curves show the analytical growth rate from equation (3.2) for a uniform buoyancy gradient equal to the peak value of the corresponding profile. (b) Peak buoyancy gradient at t_{onset} against k_{max} for all cases, coloured by particle diameter. The dash-dotted line is a least-squares fit. (c) The same gradient against k^* . The dashed line shows the analytical stability boundary from equation (3.5), evaluated at $\theta = 0$. All points lie above this boundary, consistent with $k^* \leq k_c$.

because the steepest point is the most unstable point, hence, no disturbance can experience stronger driving.

The full growth rate at RTI onset $\sigma(k, t_{\text{onset}})$ shows that growth vanishes at both ends of the wavenumber range, but for entirely different reasons (figure 5a). At small k , viscosity and diffusion are negligible and growth is limited instead by the depth of fluid that the disturbance must set in motion. A mode of wavenumber k accelerates a layer of fluid of thickness $\sim k^{-1}$ on either side of the interface, since its velocity decays as $e^{-k|z|}$ away from it. The mass set in motion therefore grows as k^{-1} while the buoyancy force driving it does not, giving $\sigma = (-\Delta b k / 2)^{1/2}$ for a two-layer interface (Smyth & Carpenter 2019, §2.2.4), where Δb is the buoyancy difference across the interface. Low wavenumber modes are thus slow because of the mass they must move rather than because they are damped.

At large k , the damping term $\nu \kappa_b K^4$ grows faster than the buoyancy forcing, so however strong the destabilizing gradient, unstable modes cease to exist beyond a finite wavenumber: σ crosses zero at k^* , and all smaller scales are stable (figure 5a). Between cases the damping is unchanged, since ν and κ_b are molecular; what differs is the buoyancy gradient. Cases with larger $\mathcal{A}$ reach onset with a weaker peak buoyancy gradient, so the destabilizing effect is smaller at every wavenumber. Because the damping grows as K^4 , it penalizes high wavenumbers far more heavily than low ones, and weakening the buoyancy forcing therefore removes the highest-wavenumber modes first. Both k^* and k_{max} are consequently smaller in cases with stronger pre-existing turbulence.

Figure 5(b) shows the peak buoyancy gradient at t_{onset} against k_{max} . Despite differences in particle size (and hence Stokes settling velocity) and pre-existing turbulence amplitude, the cases follow a single power law, $\max |\partial\langle b \rangle_{xy}/\partial z| \propto k_{\text{max}}^{4.6 \pm 0.7}$. This suggests that both D and $\mathcal{A}$ influence k_{max} through their effects on the peak buoyancy gradient. Had either acted on the instability by some other route, cases with different D or $\mathcal{A}$ would have fallen on separate curves. The $D = 30 \mu\text{m}$ cases scatter about this relation more than the other two families. The stability calculation treats the buoyancy profile as frozen at each instant and contains no settling term, whereas settling could continuously reshape the profile. The neglected evolution scales with the settling parameter W_p , which is largest for these cases, and they also reach onset earliest, when the gradient is changing most rapidly. The separation between the growth rate and the rate at which the base state evolves is therefore

smallest for these cases, so the quasi-steady approximation underlying $k_{\max}$ is least well satisfied and its correspondence with the peak gradient loosens accordingly.

Figure 5(c) shows the same peak buoyancy gradient against the neutral wavenumber k^* . The data again collapse across particle sizes and turbulence amplitudes, and lie above the analytical relation $|\partial\langle b \rangle_{xy}/\partial z| = \nu\kappa_b k^4$ (dashed), consistent with the bound $k^* \leq k_c$. The departure from the line is not uniform: it is largest for the cases with the smallest k^* , that is, the strongest pre-existing turbulence. The gap measures how closely the profile resembles a uniform gradient near its maximum, and it grows systematically with $\mathcal{A}$.

4. Transition

The stability analysis of §3 assumes that the disturbance evolves as a normal mode: a fixed spatial structure whose amplitude grows as $e^{\sigma t}$. A direct consequence is that during the linear stage the *pattern* of the perturbation should not change, only its magnitude. Figure 6 shows horizontal slices of the horizontal velocity magnitude U (first row) and vertical velocity w (second row) at each case's onset at $z = 0.05$ m for $D = 10$ μm . At low $\mathcal{A}$ the vertical velocity is organized into quasi-cellular structures of a well-defined scale, whereas at high $\mathcal{A}$ the field is dominated by elongated filaments and the cellular structure is less organized. The animation (Movie 1) shows this more clearly. In the strongly forced case ($\mathcal{A} = 0.0172$), the cellular structure itself appears broadly stationary, but the smaller-scale features surrounding it rearrange continuously, most obviously between $t = 43$ and 54 s. In the weakly forced case ($\mathcal{A} = 0.000207$), by contrast, the whole field remains essentially fixed while its amplitude intensifies toward onset ($t = 70 - 79$ s). At high $\mathcal{A}$, therefore, the flow does not maintain a persistent spatial pattern overall, even though structure at the scale of the mode itself remains identifiable.

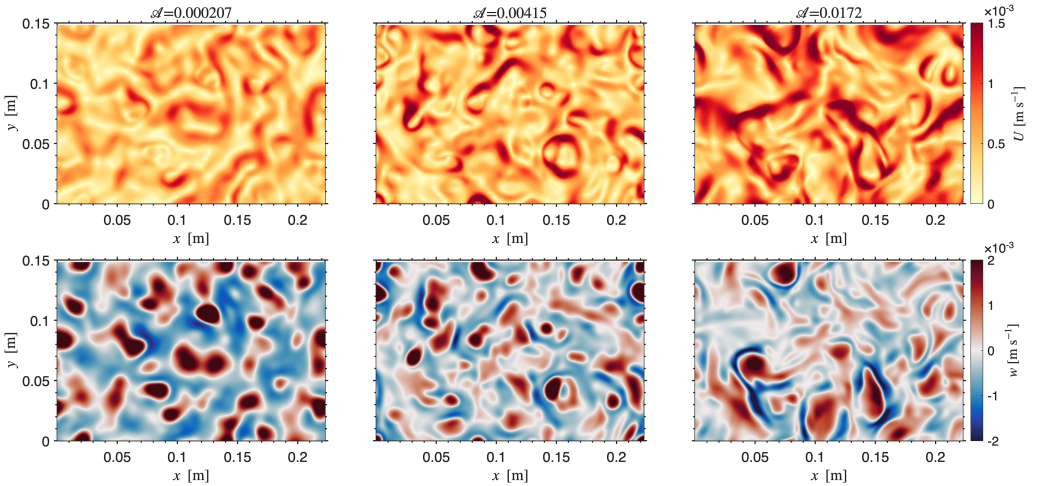


Figure 6. Horizontal ($x - y$) slices at $z = 0.05$ m for $D = 10$ μm , each taken at that case's own RTI onset time. Columns correspond to increasing pre-existing turbulence amplitude $\mathcal{A}$; the top row shows the horizontal velocity magnitude $U = (u^2 + v^2)^{1/2}$ and the bottom row the vertical velocity w .

The contrast between the sections in figure 6 can be made quantitative by asking how the kinetic energy at onset is divided between motion in the plane of the section and motion across it, and over which scales. We take the same horizontal sections as in figure 6, transform each velocity component in x and y , and sum the resulting power into annular

bands of horizontal wavenumber $k_h = (k_x^2 + k_y^2)^{1/2}$,

$$E_U(k_h) = \frac{1}{\Delta k} \sum_{\mathbf{k}_h \in \mathcal{K}(k_h)} \left(|\hat{u}|^2 + |\hat{v}|^2 \right), \quad E_w(k_h) = \frac{1}{\Delta k} \sum_{\mathbf{k}_h \in \mathcal{K}(k_h)} |\hat{w}|^2, \quad (4.1)$$

where

$$\mathcal{K}(k_h) = \left\{ \mathbf{k}_h = (k_x, k_y) : |\mathbf{k}_h| \in \left[k_h - \frac{1}{2}\Delta k, k_h + \frac{1}{2}\Delta k \right] \right\} \quad (4.2)$$

is the annular band of width Δk centred on k_h , and carets denote the two-dimensional Fourier transform over the section. E_U measures horizontal motion and E_w vertical motion, both as functions of the horizontal scale on which that motion occurs. The comparison between them is informative here because the two have different origins. The pre-existing turbulence is imposed as horizontal motion and is further constrained by the stratification, whereas the instability converts potential energy into vertical motion at the scale it selects, so E_w is the part of the spectrum that the instability builds and E_U is largely the environment it has to build in.

Early in each run the curves for different $\mathcal{A}$ are widely separated (figure 7), since at that stage the field is essentially the imposed perturbation and its amplitude fixes the energy at every scale. That separation restates the initial condition rather than telling us anything about the instability. By the onset the instability has amplified the vertical motion in every case and the two spectra have converged to comparable magnitudes. For the weakest pre-existing turbulence, E_w exceeds E_U near the wavenumber selected by the linear analysis of §3, so the vertical motion of the emerging cells is stronger than the horizontal motion surrounding them (figure 7a). This margin narrows as $\mathcal{A}$ increases, and at $\mathcal{A} = 1.72 \times 10^{-2}$ the horizontal energy has been enhanced to the point where the peak of E_w is no longer larger than E_U but comparable to it (figure 7c). Their ratio therefore measures how far the instability stands out from the flow it grows in. Where E_w exceeds E_U the motion at the selected scale is overturning, which is the signature of the instability having taken charge of the flow. Where the two are comparable the motion retains the flattened, largely horizontal character imposed at the start, and the horizontal motion surrounding the emerging cells is then as strong as their own overturning, strong enough to strain the pattern, creating filamentary structures (third column in figure 6).

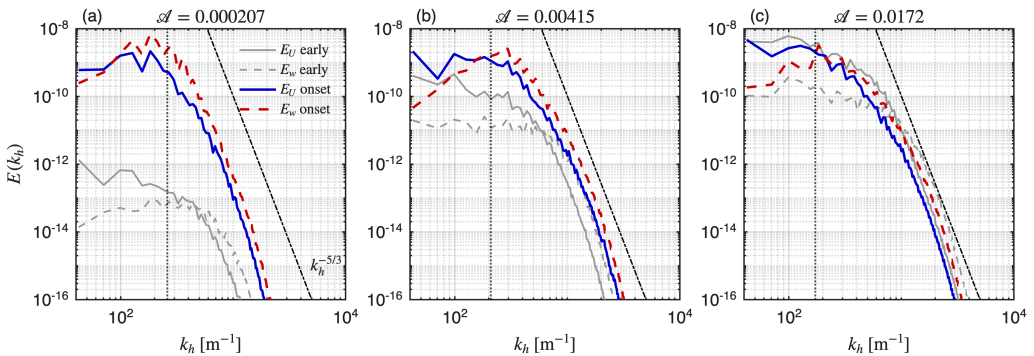


Figure 7. Radially binned horizontal and vertical kinetic energy spectra, $E_U(k_h)$ and $E_w(k_h)$ from equation (4.1), computed from horizontal sections at $z = 0.05$ m for $D = 10 \mu\text{m}$ at (a) $\mathcal{A} = 0.000207$, (b) $\mathcal{A} = 0.00415$ and (c) $\mathcal{A} = 0.0172$. Each panel shows an early time (grey) and that case's own onset (coloured), with solid curves for E_U , dashed for E_w , and the vertical line marking k_{max} from the linear analysis. Across panels the early-time levels differ according to the imposed amplitude but have converged by onset. Near k_{max} , E_w exceeds E_U at the weakest amplitude, whereas the two are comparable at the strongest.

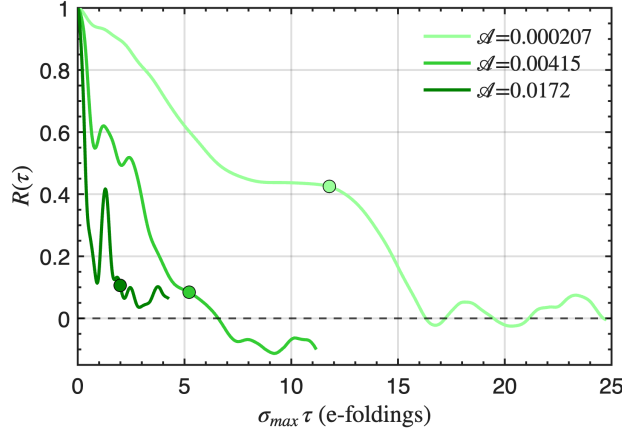


Figure 8. Temporal autocorrelation $R(\tau)$ of the vertical velocity on the horizontal plane $z = 0.05$ m, defined by equation (4.3), for three pre-existing-turbulence amplitudes at $D = 10 \mu\text{m}$, with darker shading indicating larger $\mathcal{A}$. The reference field is taken at $t_0 = 0.7t_{\text{onset}}$ and the lag τ is expressed in units of the instability e-folding time, $\sigma_{\text{max}}\tau$, so that each case is sampled over the same relative stage of its evolution. Normalization by the r.m.s. removes the exponential amplification, so $R = 1$ indicates an unchanged spatial pattern and $R \rightarrow 0$ an unrelated one; a growing normal mode would give $R \approx 1$ at all lags. Filled circles mark t_{onset} , the end of each averaging window. R decays more rapidly as $\mathcal{A}$ increases, with mean values $\bar{R} = 0.63, 0.42$ and 0.34 .

To quantify, we correlate the vertical velocity on a horizontal plane with itself at increasing time lag,

$$R(\tau) = \frac{\langle w'(\mathbf{x}_h, t_0) w'(\mathbf{x}_h, t_0 + \tau) \rangle_{xy}}{\left[\langle w'^2(t_0) \rangle_{xy} \langle w'^2(t_0 + \tau) \rangle_{xy} \right]^{1/2}}, \quad (4.3)$$

where the horizontal average is taken at $z = 0.05$ m. The covariance is normalized by the root-mean-square so that the exponential amplification cancels and R responds only to changes in spatial structure: $R = 1$ if the pattern is unchanged and the disturbance has simply grown, and $R \rightarrow 0$ once it has been rearranged into an unrelated configuration. A growing normal mode therefore gives R closer to 1 at all lags by construction, and any decay measures the rate at which the flow departs from that description. The reference field is taken at $t_0 = 0.7t_{\text{onset}}$ and the lag followed to t_{onset} , so that every case is sampled over the same relative stage of its evolution and the comparison is unaffected by the different absolute times at which onset occurs. The lag is reported in units of the instability's own growth time, $\sigma_{\text{max}}\tau$, since the physically relevant question is whether the disturbance retains its structure long enough to amplify appreciably. Each case is summarized by the mean correlation $\bar{R}$ over the interval from $0.7t_{\text{onset}}$ to t_{onset} with $D = 10 \mu\text{m}$. We note that this comparison is conservative: because σ_{max} is larger and t_{onset} later at low $\mathcal{A}$, the same relative interval spans approximately 12 e-foldings at $\mathcal{A} = 2.07 \times 10^{-4}$ but only 1.8 at $\mathcal{A} = 1.72 \times 10^{-2}$ (figure 8). Because a longer window allows more time for decorrelation, the unequal window lengths penalize the smaller $\mathcal{A}$ cases most. Nonetheless, figure 8 shows that R decays progressively more rapidly as $\mathcal{A}$ increases, with $\bar{R} = 0.63, 0.42$ and 0.34 for $\mathcal{A} = 2.07 \times 10^{-4}, 4.15 \times 10^{-3}$ and 1.72×10^{-2} , respectively. This confirms the visual impression shown by figure 6. In the most strongly perturbed case the disturbance loses coherence within roughly one e-folding, so the pattern the instability is attempting to establish is partially destroyed before it can amplify. At high $\mathcal{A}$ the disturbance is therefore rearranged before it can be regarded as a growing normal mode, and the assumption

underlying the stability analysis breaks down. The instability is then not only slower to grow but also less able to establish the organized structure through which sediment is transported, and we examine the consequences for the settling flux in §6.

5. Flow evolution

This section describes how the full three-dimensional flow evolves from the onset of the instability until the sediment first reaches the bottom boundary, beyond which the dynamics are no longer of interest. Figure 9 shows cross-sections of the concentration field for three pre-existing-turbulence amplitudes $\mathcal{A}$, each captured at three characteristic times: onset (t_{onset} , top row), the moment the RTI plumes cross $z = 0$ (middle row), and the time the sediment reaches the bottom, t_f (bottom row).

At onset, the interface is already visibly less coherent for larger $\mathcal{A}$: the largest-amplitude case (figure 9g) shows small-scale structure that is absent in the weakly forced case (figure 9a), consistent with the loss of a well-defined normal mode described in §4. Onset also occurs sooner: as in figure 4a, the strongest pre-existing turbulence reduces t_{onset} to 67% of its value at the weakest amplitude. By the time the plumes cross $z = 0$ (middle row), the contrast has sharpened further: the plumes remain sharp and coherent at low $\mathcal{A}$

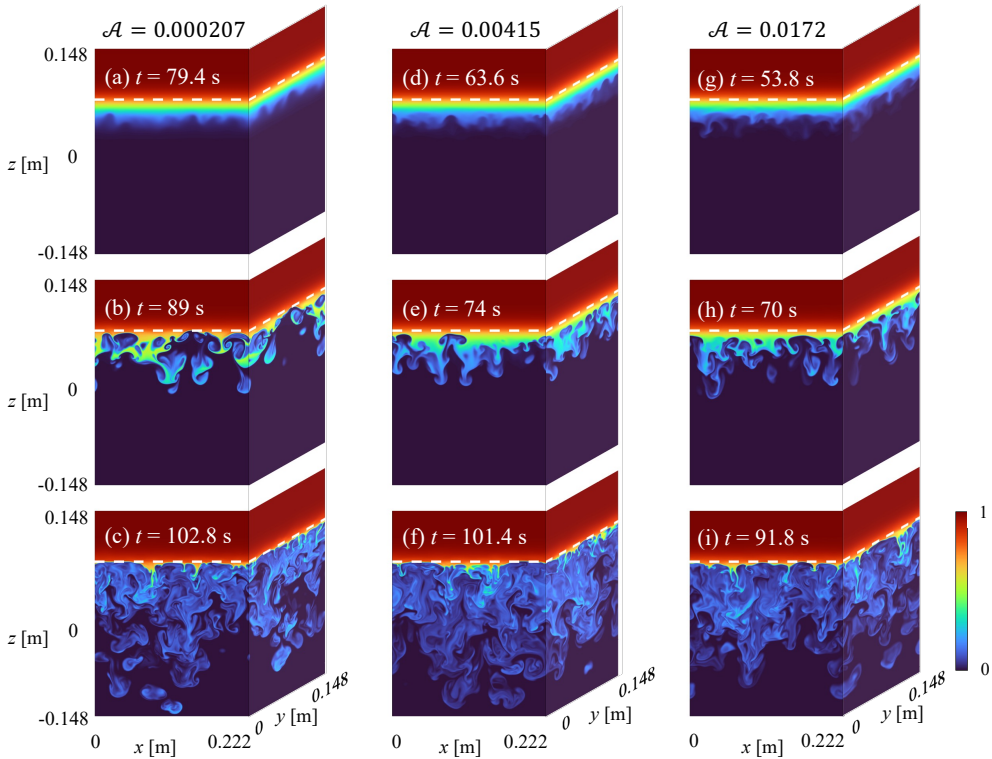


Figure 9. Cross-sections through the 3-D concentration field at successive times for three pre-existing turbulence amplitudes: (a-c) $\mathcal{A} = 0.000207$, (d-f) $\mathcal{A} = 0.00415$, (g-i) $\mathcal{A} = 0.0172$. Concentration is normalized by Δc . Rows correspond to three characteristic times: the top row shows t_{onset} ; the middle row shows the time when RTI is fully developed and crosses $z = 0$; and the bottom row shows $t = t_f$, the time at which the sediment first reaches the bottom boundary. The particle size D is $10 \mu\text{m}$, corresponds to a $W_{st} = 0.00012 \text{ m s}^{-1}$. White dashed line is the initial sediment interface.

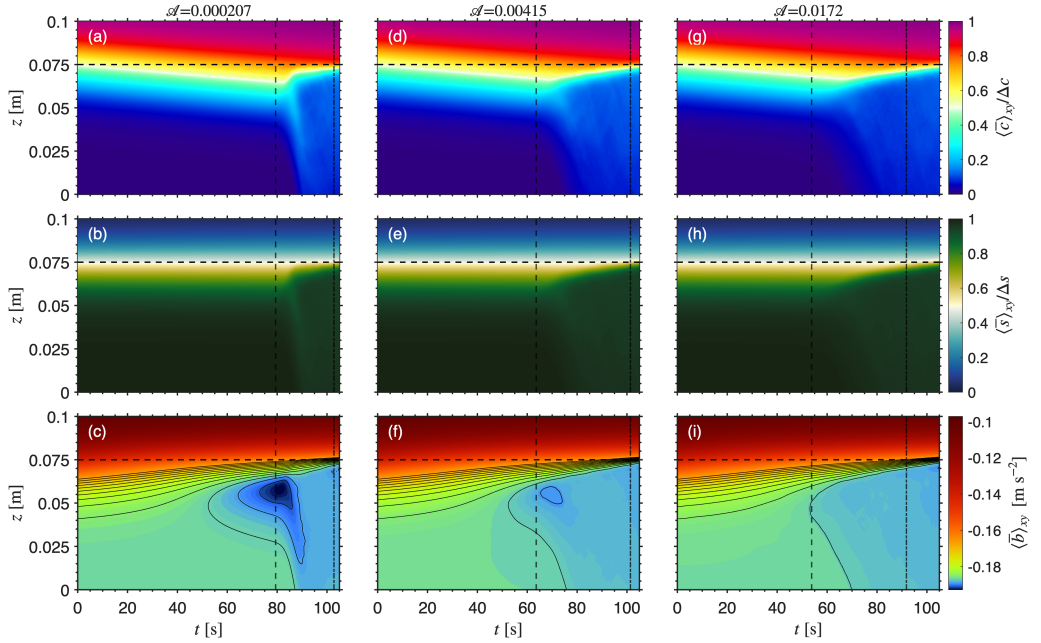


Figure 10. Hovmöller diagrams of the horizontally averaged (a,d,g) normalized sediment concentration, (b,e,h) normalized salinity, and (c,f,i) buoyancy, as functions of height z and time t , for the same three pre-existing-turbulence amplitudes as in figure 9: (a-c) $\mathcal{A} = 0.000207$, (d-f) $\mathcal{A} = 0.00415$, (g-i) $\mathcal{A} = 0.0172$. The horizontal dashed line marks the initial location of the sharp interface, $z = 0.075$ m. In each panel, the two vertical lines mark t_{onset} (dashed) and t_f (dot-dashed), corresponding to the RTI onset time and the time at which sediment first reaches the bottom boundary, respectively. Black contours in (c,f,i) show $\langle \bar{b} \rangle_{xy}$ from -0.1929 to -0.165 m s^{-2} in increments of 0.002 m s^{-2} , tracing the descent of the buoyancy interface; the closed contour marks a local buoyancy minimum, where the nose forms.

(figure 9b) but are visibly smeared, with weaker tracer gradients, at high $\mathcal{A}$ (figure 9h). By t_f (figure 9c,f,i), however, this contrast has largely disappeared: coherent plume structure is no longer evident near the interface (white dashed line) in any case, sediment instead reaching the bottom through scattered, sporadic filaments. Despite this apparent convergence in structure, the time required to get there is not the same: the interval from onset to t_f is 23 s at $\mathcal{A} = 0.000207$ but 38 s at $\mathcal{A} = 0.0172$, so the weakest-turbulence case needs only 60% as long (table 2). Pre-existing turbulence therefore has two opposing effects on timing: it advances the onset of the instability but slows the subsequent descent of sediment to the bottom. We quantify the transport consequences of this trade-off directly in §6.

To track this evolution continuously rather than at three snapshots, figure 10 shows Hovmöller diagrams of the horizontally averaged concentration, salinity, and buoyancy fields. Before onset, the sediment interface deepens steadily while the salinity interface remains essentially fixed, so the two separate into a distinct nose (white interface in figure 10a,d,g). The leading edge of this first front descends at a slope that is roughly constant in time and changes little with $\mathcal{A}$, though the front persists for a shorter time, breaking down at smaller t_{onset} , as $\mathcal{A}$ increases. Once RTI develops, the front no longer follows this steady linear descent, and a clear contrast between amplitudes emerges: at low $\mathcal{A}$ (figure 10a) it steepens sharply, plunging toward the bottom as coherent plumes carry sediment down directly, whereas at high $\mathcal{A}$ (figure 10g) the descent instead becomes flatter and smears, reflecting the weaker, less organized plumes described in §4. The buoyancy field (figure 10c,f,i) makes the fate of the nose explicit: its contours pinch off into a closed,

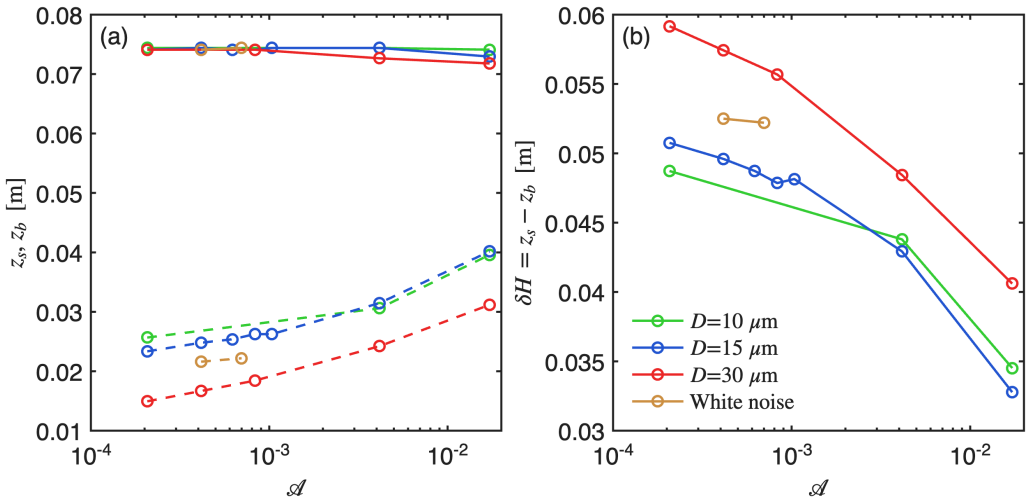


Figure 11. (a) Vertical location of the sharp salinity interface z_s (solid) and of the buoyancy threshold z_b (dashed) as a function of pre-existing turbulence amplitude $\mathcal{A}$, for particle diameters $D=10 \mu\text{m}$ (green), $D=15 \mu\text{m}$ (blue) and $D=30 \mu\text{m}$ (red), and for the white-noise initial condition (brown). (b) Corresponding nose height $\delta H = z_s - z_b$ versus $\mathcal{A}$. The nose height decreases with increasing turbulence amplitude for all cases, driven primarily by the increase in z_b rather than by changes in z_s .

isolated pocket of low buoyancy that marks the nose detaching from the retreating salinity interface as the plunging front overtakes it. This pocket is prominent at low $\mathcal{A}$ (figure 10c), noticeably smaller at intermediate $\mathcal{A}$ (figure 10f), and much less discernible at the highest amplitude (figure 10i).

Figure 10 illustrates this suppression of the nose qualitatively for a single particle size, $D = 10 \mu\text{m}$; we now quantify it across the full set of amplitudes, particle sizes, and the white-noise perturbation. The nose height δH , which serves as a useful indicator of RTI strength, is reduced by pre-existing turbulence. We define $\delta H = z_s - z_b$, where z_s denotes the location of the sharpest salinity interface and z_b denotes the location at which buoyancy first falls below the bottom-layer buoyancy by 0.2% (illustrated in figure 1). The trend in nose height across cases is not sensitive to the choice of cutoff percentage; however, because the nose region is considerably smaller for large cases of pre-existing turbulence (e.g. $\mathcal{A} = 0.0172$), a larger cutoff percentage would fail to capture the lower extent of the nose.

The nose height decreases with increasing pre-existing-turbulence amplitude for all three particle sizes (figure 11b), by a comparable amount overall. At low $\mathcal{A}$, δH increases with particle size, but $D = 10$ and $D = 15$ nearly coincide over an intermediate range of $\mathcal{A}$ and the strongest amplitude. The position of the salinity interface, z_s , is nearly independent of $\mathcal{A}$ for all particle sizes (solid curves in figure 11a), so z_b dominates the trend in δH .

A direct comparison is available at $D = 15 \mu\text{m}$ between a pre-existing field with a $k^{-5/3}$ spectrum (blue) and one with a white spectrum (brown); both are randomly phased, so the two differ only in how energy is distributed across horizontal scales. Since they differ in energy at equal $\mathcal{A}$, case 5 is compared with white noise matched both in nominal amplitude (case 12) and in total kinetic energy (case 11); under both matchings the nose sits about 0.003 m higher and onset is 4–5 s later than the pre-existing case. The $k^{-5/3}$ field is therefore more effective than white noise at bringing the instability on early, and the shallower nose follows from that earlier onset.

6. Sediment flux and effective settling speed

Sections 3 and 4 established what pre-existing turbulence does to the instability itself. A larger amplitude $\mathcal{A}$ brings the instability on earlier, at a stage when the settling nose is less developed, so the buoyancy gradient available to drive the instability is weaker and both the growth rate and the fastest-growing wavenumber are reduced. The emerging pattern is also less coherent, and the nose that the instability grows on is shallower. None of this yet says how much sediment is actually moved. The instability matters here because it transports sediment downward faster than a particle would fall on its own, so the question this section answers is whether the systematic weakening of the instability translates into a systematic reduction in transport.

6.1. Vertical sediment flux and turbulent diffusivity

The vertical sediment flux is $\langle w' \hat{c}' \rangle_{xy}$, where $\hat{c} = c/\Delta c$ is the normalized sediment concentration and primes denote departures from the horizontal average. Its evolution over the full depth is shown in the upper row of figure 12. Before onset the flux is indistinguishable from zero. At onset a strongly negative band appears at the sediment interface and then descends with the plumes, and it is this band that carries the sediment delivered to depth. Comparing across columns, the flux weakens monotonically as $\mathcal{A}$ increases (figure 12a-c). Averaging over the interval from onset to the time the sediment

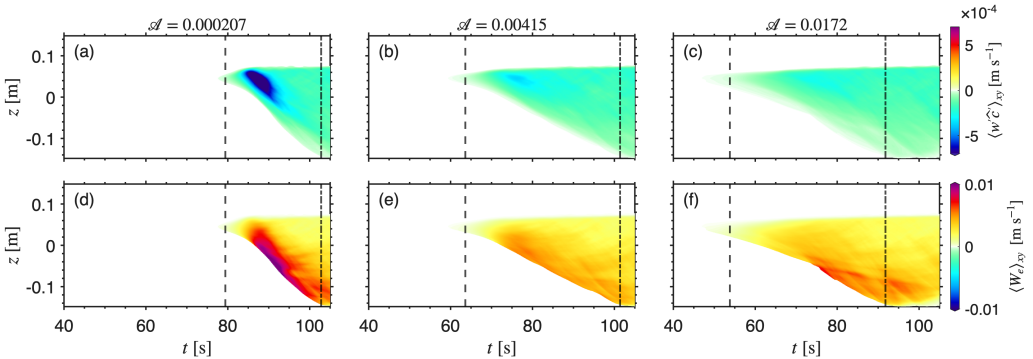


Figure 12. Vertical sediment flux and effective settling velocity for $D = 10 \mu\text{m}$, with the pre-existing turbulence amplitude increasing from left to right. (a–c) Horizontally averaged sediment flux $\langle w' \hat{c}' \rangle_{xy}$ as a function of depth and time. (d–f) Effective settling velocity $\langle W_e \rangle_{xy} = -\langle w' c' \rangle_{xy} / \langle c \rangle_{xy}$ over the same window. Dashed and dot-dashed vertical lines mark the onset of the instability and the time at which sediment first reaches the bottom, respectively. Levels at which $|\langle c \rangle_{xy}| / \Delta c < 10^{-3}$ are blanked in the lower row, and only the strongly negative portion of the sediment flux is retained in the upper row.

first reaches the bottom and over the depth of the layer (figure 13a) reveals that between the weakest and strongest pre-existing turbulence the volume-averaged flux differs by roughly a factor of two, and the ordering with $\mathcal{A}$ is the same as the ordering of the linear growth rate in §3. The consequence of a weaker, less coherent instability is therefore a weaker flux.

The flux is measured directly and requires no closure assumption, but an alternative way to measure it is to adopt the flux–gradient description and define turbulent diffusivities for sediment and for salt (Burns & Meiburg 2015),

$$K_c = \frac{-\overline{\langle w' c' \rangle_{xy}}}{d\langle c \rangle_{xy}/dz}, \quad K_S = \frac{-\overline{\langle w' S' \rangle_{xy}}}{d\langle S \rangle_{xy}/dz}. \quad (6.1)$$

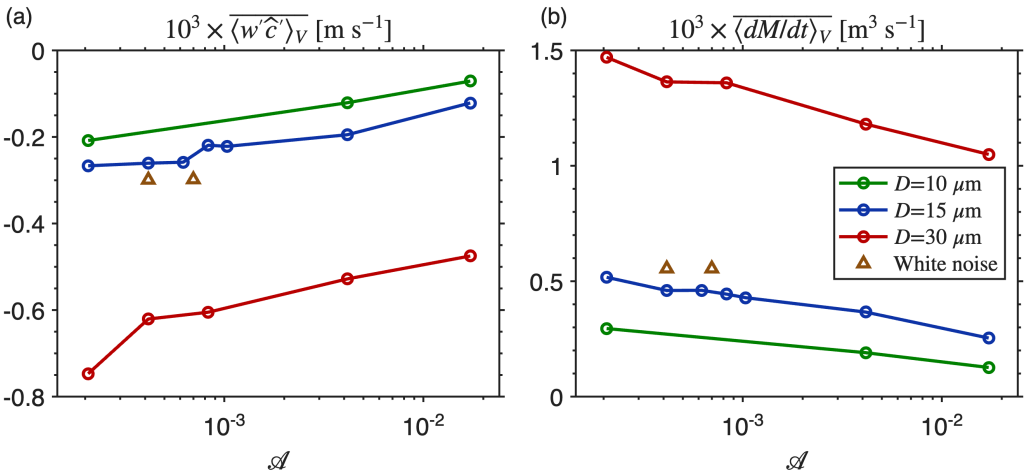


Figure 13. Quantities averaged over the interval from the onset of the instability to the time at which sediment first reaches the bottom, against pre-existing turbulence amplitude $\mathcal{A}$. (a) Volume-averaged sediment flux and (b) volume-averaged mass rate dM/dt for the layer beneath z_{RT} . Lines denote the three particle-size families and triangles the white-noise initial condition. The magnitude of both quantities decreases monotonically with $\mathcal{A}$, by approximately a factor of two over the range examined. Note that the sediment flux in (a) is negative.

The overbar denotes an average over the vertical band $z_c \pm \lambda_{\max}$, where z_c is the level of steepest mean gradient at that instant and $\lambda_{\max} = 2\pi/k_{\max}$ is the wavelength of the fastest-growing mode taken from the linear stability analysis of §3. The gradient in the denominator is evaluated at z_c itself. Following the level of steepest gradient rather than a fixed height means that the diffusivity tracks the descending interface, which is where the instability is strongest and where almost all of the flux resides. Using $\lambda_{\max}$ to set the width of the band ties the measurement to the scale the instability actually selects, so the same physical region is sampled in every case even though that scale itself varies with $\mathcal{A}$.

Figure 14 shows that before onset both diffusivities lie below the molecular value, $K_c/\kappa_c < 1$ and $K_S/\kappa_S < 1$, because the flux at that stage is negligible. The pre-existing turbulence, however energetic, is not organized in a way that carries sediment across the interface. At onset (dots) the diffusivities rise abruptly by one to two orders of magnitude as the plumes form. They peak shortly afterwards and then decay, and within a particle-size family the late-time values converge. This is because both the mean profile and the flux at the interface become nearly independent of $\mathcal{A}$. The nonlinear phase therefore keeps little memory of the initial perturbation, since its only lasting influence is on the time at which the instability starts and hence on how far the settling nose had developed by that moment. What $\mathcal{A}$ sets is the peak the diffusivity reaches, not the state it relaxes to. The peak diffusivity $\max(K_c)/\kappa_c$ decreases nearly monotonically with $\mathcal{A}$ (figure 14c), and the same is true for every particle size. This is the same ordering as the flux in figure 13a. At the weakest pre-existing turbulence the instability mixes sediment across the interface some two orders of magnitude faster than molecular diffusion, and pre-existing turbulence erodes that advantage systematically.

Comparing the two diffusivities is informative because the same velocity field carries both scalars and, with $\kappa_c = \kappa_S$, two passive scalars would give $K_c = K_S$ exactly. Neither scalar is passive here. Both scalars acquire fluctuations in the same way, through displacement of their mean profiles, so the difference between them lies not in how the fluctuation is created but in what becomes of the displaced parcel. A parcel carried downward is fresher than its surroundings, and the buoyancy anomaly that follows is

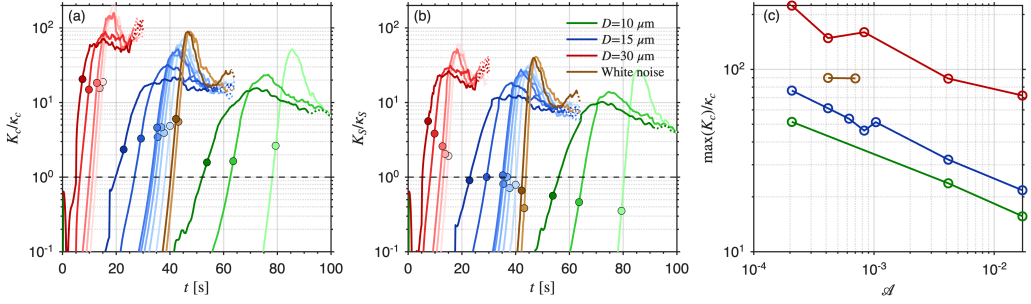


Figure 14. Turbulent diffusivities of sediment and salt, defined by equation (6.1) with the flux averaged over $z_c \pm \lambda_{\max}$ and the gradient evaluated at z_c , the level of the steepest mean gradient. (a) Sediment diffusivity K_c normalized by the molecular diffusivity κ_c , and (b) salinity diffusivity K_s , normalized by κ_s , as functions of time. Colours follow the case groups of figure 3, with darker shades indicating larger $\mathcal{A}$; filled circles mark the onset of the instability, the dashed horizontal line marks $K/\kappa = 1$, and the portion of each curve after the sediment reaches the bottom is drawn in dots. (c) Peak sediment diffusivity $\max(K_c)/\kappa_c$ over the interval from onset to bottom arrival, against $\mathcal{A}$.

restoring, so the parcel tends to return and some of the salinity flux can be undone. Its sediment anomaly acts the other way, making the parcel heavier than its surroundings so that the displacement grows, and the sediment is delivered rather than returned. The excess of K_c over K_s is a measure of that difference, and therefore of the part of the transport that belongs to RTI. Figure 14(a,b) shows that the inequality can be seen at all particle sizes, especially for smaller particles, where K_s barely exceeds κ_s at onset while K_c is already well above κ_c . It is also a useful marker of when the instability takes over. Before onset the only motions present are those of the imposed turbulence, which has no reason to prefer one scalar over the other, so the separation between K_c and K_s that opens at onset identifies the point at which the transport ceases to be that of the initial field and becomes that of the instability.

6.2. Effective settling velocity

For applications, the quantity of interest is the speed at which sediment descends, and the conventional measure (Yu *et al.* 2014; Burns & Meiburg 2015; Yang *et al.* 2025) is the effective settling velocity

$$\langle W_e \rangle_{xy}(z, t) = -\frac{\langle w'c' \rangle_{xy}}{\langle c \rangle_{xy}}, \quad (6.2)$$

defined as the ratio of the horizontally averaged sediment flux to the horizontally averaged sediment concentration. The effective settling velocity is largest near the leading edge of the plunging plume (figure 12(d–f)), as expected since the fastest downward motions occur there. Comparing the weakest and strongest pre-existing turbulence cases, $\mathcal{A} = 2.07 \times 10^{-4}$ and $\mathcal{A} = 1.72 \times 10^{-2}$, respectively, both the flux and the settling speed are reduced, so at first sight $\langle W_e \rangle_{xy}$ appears to inherit the ordering of the flux. Once the intermediate case $\mathcal{A} = 4.15 \times 10^{-3}$ is included, however, this ordering breaks down. The flux remains monotonic in $\mathcal{A}$, consistent with the linear results, yet $\langle W_e \rangle_{xy}$ for $\mathcal{A} = 1.72 \times 10^{-2}$ is larger than that for $\mathcal{A} = 4.15 \times 10^{-3}$. The reason is a limitation of the definition rather than a physical reversal. Equation (6.2) divides by $\langle c \rangle_{xy}$, and near the lower edge of the front the mean concentration approaches zero. A small flux divided by an even smaller concentration yields a large velocity, so there are regions in figure 12 where $\langle W_e \rangle_{xy}$ is large while the flux there is negligible. The velocity is not incorrect, but it applies to a vanishingly small fraction of the sediment. When reported as a single number it overstates how fast

the material as a whole is moving, and it is exactly this that produces the apparent non-monotonicity. The remedy is to weight by the amount of sediment present. We therefore define a bulk effective settling velocity by integrating the flux and the concentration over the layer from the bottom to z_{RT} ,

$$\overline{W_e}^{\text{bulk}}(t) = -\frac{\int_{-L_z/2}^{z_{RT}} \langle w' c' \rangle_{xy} dz}{\int_{-L_z/2}^{z_{RT}} \langle c \rangle_{xy} dz} = \frac{\int_{-L_z/2}^{z_{RT}} \langle W_e \rangle_{xy} \langle c \rangle_{xy} dz}{\int_{-L_z/2}^{z_{RT}} \langle c \rangle_{xy} dz}, \quad (6.3)$$

where z_{RT} is the depth of the sediment interface at the time of onset. The second form makes the interpretation explicit. Equation (6.3) is a concentration-weighted average of $\langle W_e \rangle_{xy}$, so that each level contributes in proportion to the amount of sediment it contains, while near-empty regions that contaminate the pointwise definition carry negligible weight.

Normalizing by the Stokes settling velocity gives the enhancement factor

$$\chi^{\text{bulk}} = \overline{W_e}^{\text{bulk}} / W_{st}. \quad (6.4)$$

For the larger particles the bulk of the sediment settles at essentially the same speed regardless of $\mathcal{A}$ (figure 15a). For $D = 10 \mu\text{m}$ the enhancement factor varies by more than a factor of two across the range of $\mathcal{A}$ examined. The sensitivity to pre-existing turbulence therefore increases as particle size decreases. For large particles, the Stokes settling velocity is already large and the instability contributes only a modest enhancement, so weakening the instability has relatively little effect on the total settling speed. For small particles, the instability accounts for most of the downward transport, so weakening the instability has a correspondingly greater effect.

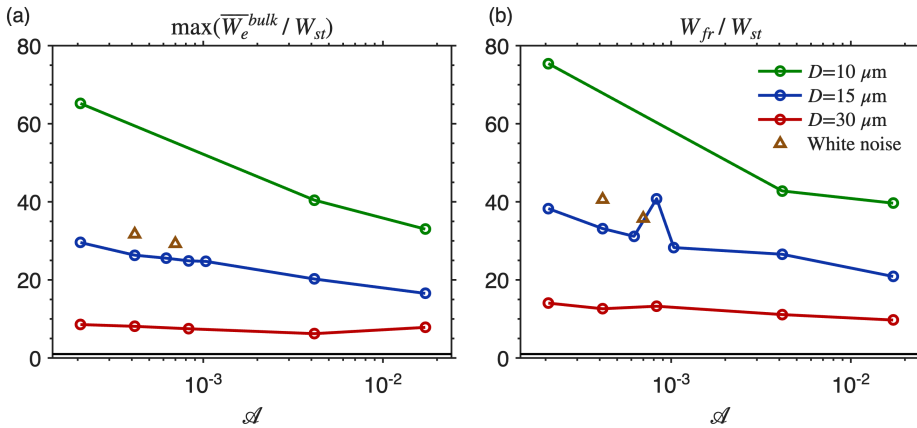


Figure 15. Settling speed relative to the Stokes velocity, against pre-existing turbulence amplitude. (a) Enhancement factor $\max(\overline{W_e}^{\text{bulk}}) / W_{st}$ obtained from the bulk definition in equation (6.3), and (b) speed inferred from the slope of the frontal trajectories in figure 16, W_{fr} / W_{st} . The solid horizontal line marks pure Stokes settling. The two estimates differ in magnitude, the front being a leading-edge measure and the bulk velocity a mass-weighted mean, but agree in trend. Both are relatively insensitive to $\mathcal{A}$ for the largest particles and vary by more than a factor of two for $D = 10 \mu\text{m}$

An independent estimate comes from the descending front. For each depth we record the first time at which the magnitude of the horizontally averaged sediment flux there exceeds 1% of its own time-averaged magnitude, and read as a function of depth this arrival time defines the frontal position $z_{fr}(t)$. Figure 16 shows the front descending far more steeply than pure Stokes settling (dashed lines) in every case. A linear fit to each

trajectory gives its slope, which we take as the frontal settling velocity, W_{fr} , and W_{fr}/W_{st} is of order ten (figure 15b). That is slightly larger than the bulk enhancement factor, since the front is the leading edge set by the fastest plumes whereas $\overline{W_e}^{bulk}$ is a mass-weighted mean over all the sediment. The two nonetheless agree where it matters, both being nearly insensitive to $\mathcal{A}$ for the largest particles and both varying by close to a factor of two for $D = 10 \mu\text{m}$. Within a family, the dependence of W_{fr} on $\mathcal{A}$ is not monotonic, reaching a maximum at intermediate amplitude for $D = 15 \mu\text{m}$, and the criterion is the likely cause. Since each depth is referenced to its own time average, the threshold falls wherever the flux is weak, and the crossing time there records a fluctuation rather than the arrival of the plume. Differences in fitted slope within a family could therefore be misleading. A criterion of this kind is sensitive to the choice of the initial perturbation (appendix A): repeating the $\mathcal{A} = 8.3 \times 10^{-4}$ case with a different set of random phases moves W_{fr} by more than half the range covered by the whole family, so the apparent maximum falls within the scatter between individual realizations. We use the front to confirm the magnitude of the enhancement and its dependence on particle size, and rely on equation (6.3) for the dependence on $\mathcal{A}$.

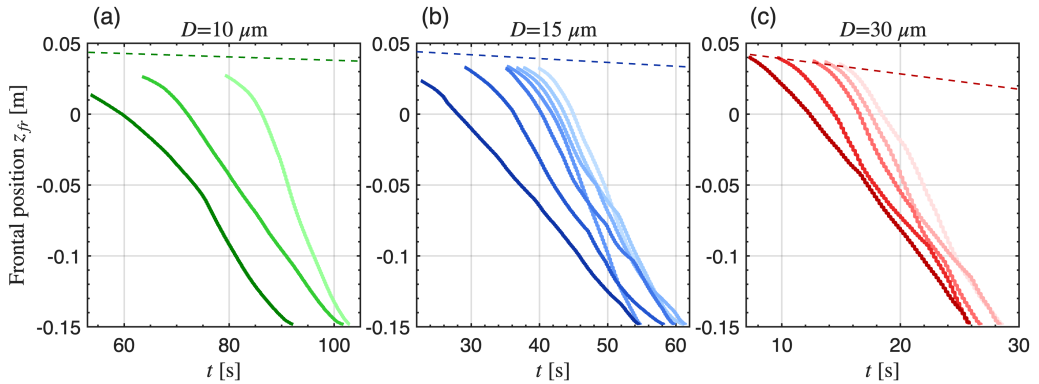


Figure 16. Vertical position of the descending front for (a) $D = 10 \mu\text{m}$, (b) $D = 15 \mu\text{m}$ and (c) $D = 30 \mu\text{m}$, with darker shades indicating larger $\mathcal{A}$. The front at a given depth is defined as the first time at which $|\langle w' \hat{c} \rangle_{xy}|$ exceeds 1% of its temporal average at that depth, and only the portion below the depth reached at onset is shown. Dashed lines show the trajectory of a particle settling at the Stokes velocity W_{st} for each particle-size family. The front descends much faster than Stokes settling in every case.

6.3. Delivered mass and mass rate

609

610 Another quantity of interest is how much sediment is transported below the interface and
 611 at what rate. We therefore integrate the mean concentration over the layer beneath z_{RT} to
 612 obtain the delivered mass

$$M(t) = \int_{-L_z/2}^{z_{RT}} \iint \langle \hat{c} \rangle_{xy} dx dy dz. \quad (6.5)$$

613 M measures the total sediment that has been carried below the depth at which the instability
 614 began, and dM/dt the rate at which it arrives there.

615 Figure 17 shows $M(t)$ and dM/dt . Both are flat before onset and rise sharply afterwards,
 616 each case starting at its own onset time, so the curves are displaced along the time axis in
 617 the order established in §3. By the time the sediment reaches the floor, M has reached much
 618 the same value for every case within a particle-size family and shows no clear ordering in
 619 $\mathcal{A}$ (figure 17a). The mass rate behaves differently. Its time series separates by amplitude

throughout the evolution, the lighter curves lying above the darker ones (figure 17b), so the ordering is a feature of the whole descent rather than of a particular instant. Averaging over the interval from onset to bottom arrival and plotting against $\mathcal{A}$ (figure 13b) puts the reduction at about a factor of two across the range examined.

Pre-existing turbulence therefore governs when delivery begins and how fast it proceeds, but not how much has been delivered once the front has crossed the layer. Writing the delivered mass as the delivery rate multiplied by the transit time, and the transit time as the layer depth divided by the frontal speed, gives $M \approx (\overline{dM/dt})H/W_{fr}$. Pre-existing turbulence reduces both the delivery rate and the frontal speed by similar factors, so their ratio is nearly unchanged and M is fixed by the depth of the layer. Stronger pre-existing turbulence weakens the instability and moves the same sediment more slowly rather than moving less of it.

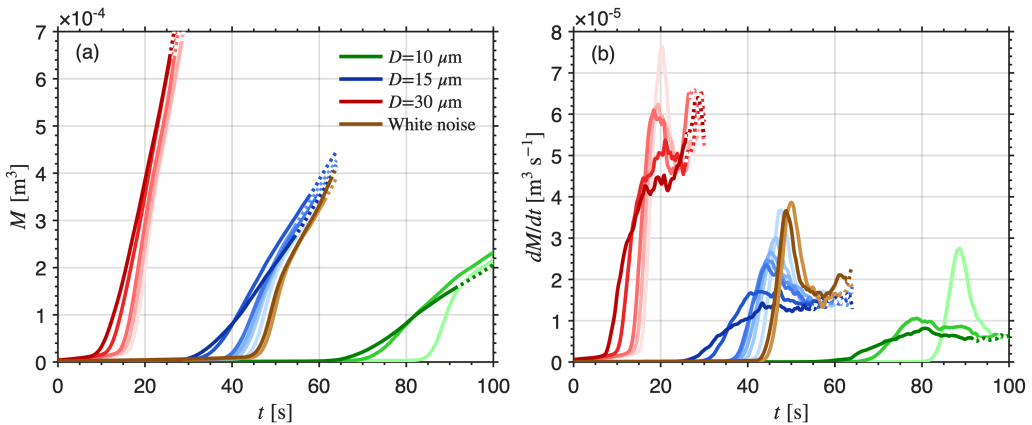


Figure 17. (a) Sediment mass M delivered below the depth z_{RT} of the interface at onset, defined by equation (6.5), and (b) the corresponding delivery rate dM/dt . Colours follow the case groups, with darker shades indicating larger $\mathcal{A}$, and the portion of each curve after the sediment reaches the bottom is shown as dashed lines. By the time the sediment reaches the bottom, M has reached a similar value within each particle-size family with no clear ordering in $\mathcal{A}$, whereas the peak of dM/dt decreases monotonically with $\mathcal{A}$.

7. Discussion and conclusion

Particles settling across a density interface accumulate beneath it, producing a layer that is denser than the fluid below and driving an instability that carries sediment downward far faster than the particles would settle individually (Hoyal *et al.* 1999; Blanchette & Bush 2005; Davarpanah Jazi & Wells 2016; Hill & Lintern 2021). Linear stability analyses and direct numerical simulations have explored the mechanisms of convective sedimentation and quantified the resulting sediment fluxes (Yu *et al.* 2013, 2014; Burns & Meiburg 2012, 2015; Shao *et al.* 2017). A common feature of these studies is a quiescent initial state in which the interface is seeded with small-amplitude perturbations. Most natural systems in which such sediment-laden layers occur, however, are not quiescent. Particles are released into or carried through a turbulent surface layer whose dissipation rate spans several orders of magnitude depending on wind, wave and buoyancy forcing (Terray *et al.* 1996; Thorpe *et al.* 2003; Esters *et al.* 2018; Zeiden *et al.* 2023). For stratified shear instabilities, pre-existing turbulence has been shown to alter the timing and the transition of the instability, as well as the resulting turbulent flow (Kaminski & Smyth 2019). We focus on particle

diameters of 10–30 μm , a range in which collective settling-driven convection, rather than individual settling, controls vertical transport.

7.1. Summary

In the linear phase, pre-existing turbulence acts through the onset time. A larger amplitude $\mathcal{A}$ provides a larger initial disturbance, so the perturbation reaches a detectable amplitude sooner, and it does so while the base state is still evolving. Sediment settling across the salinity interface continually builds the dense nose on which the instability feeds, so earlier onset occurs on a less developed nose with a weaker buoyancy gradient. No other appreciable process contributes to the formation of this buoyancy gradient. Its evolution is therefore governed by the tracer settling rather than by the perturbation, and within a particle-size family, $\max |\partial b / \partial z|$ follows nearly the same trajectory in time regardless of the value of $\mathcal{A}$. Cases differ in their buoyancy gradient at onset not because the gradient evolves differently, but because each samples a common trajectory at a different time.

The same argument accounts for the shift of the fastest-growing and marginal wavenumbers to lower values. The buoyancy forcing that drives the instability is independent of wavenumber, whereas viscous and diffusive damping increase as K^4 , so the unstable band terminates where the two effects balance. The damping is the same in every case and only the buoyancy forcing changes. Weakening the forcing moves this balance to lower wavenumber, which narrows the unstable band and reduces the fastest growing wavenumber. The analytical estimate reproduces this trend, but it is derived for a uniform gradient whereas ours varies with depth, and evaluating it with the largest gradient anywhere in the profile makes it an upper bound on the observed cut-off rather than a direct prediction. It is used here as a reference scale for the trend and not as a quantitative prediction.

The shift to lower wavenumbers may seem counterintuitive with the flow itself, since larger $\mathcal{A}$ produces qualitatively more small-scale motion (figure 6) while the instability selects a larger scale. The two describe different fields. Both $k_{\max}$ and k^* belong to the linear problem posed by the mean buoyancy profile, in which the turbulence has no direct impact, whereas the fine structure seen in the horizontal sections is the decaying pre-existing turbulence together with the tracer filaments it draws out of the interface, neither of which shows evidence of being a growing mode. The small-scale motion instead strains the growing mode. The pattern correlation $R(\tau)$ measures how long a structure retains its shape once its amplitude growth is normalized away. A mode that loses its shape sooner spends fewer e-foldings growing as a coherent structure, so at larger $\mathcal{A}$ the instability enters in a less organized state.

In the nonlinear phase the consequences for transport are systematic. The vertical sediment flux, the turbulent diffusivities, the bulk effective settling velocity, and the mass rate all decrease as $\mathcal{A}$ increases over the range examined, and the frontal speed independently confirms both the magnitude of the enhancement over Stokes settling and its dependence on particle size. Reaching that conclusion requires care in defining settling velocity, since the conventional measure divides by a mean concentration that approaches zero at the lower edge of the front. A negligible flux divided by an even smaller concentration yields a large velocity, so the definition reports rapid settling in regions that contain almost no sediment. Weighting by the amount of sediment therefore provides a more representative estimation of the settling velocity. One further result constrains the interpretation. The total sediment transported by the time the front reaches the bottom boundary is nearly the same for every case within a particle size family and shows no clear dependence on $\mathcal{A}$, even though the rate at which mass is delivered varies nearly monotonically. This invariance is specific to the present small domain, in which the front reaches the bottom while the plumes remain vigorous. Over the depth of a real mixed

layer $O(10\text{--}100\text{ m})$ the descent would take far longer, and particle dissolution, plume entrainment and decay, and continuous replenishment of the turbulent field would all have time to act. Whether the delivered sediment mass remains independent of $\mathcal{A}$ over such depths is an open question.

Sensitivity to pre-existing turbulence increases as particle size decreases, whether measured by the onset time, bulk settling velocity or frontal speed. The fact that the settling enhancement produced by RTI is largest for fine particles is already established (Hoyal *et al.* 1999; Yu *et al.* 2014; Burns & Meiburg 2015). This is why the mechanism matters in settings dominated by fine material, such as estuaries and mineral-based ocean alkalinity enhancement, where individual settling is far too slow to account for the rates at which material is observed to leave the surface layer (Lu *et al.* 2022; Yang & Timmermans 2024; Qiu *et al.* 2026). Pre-existing turbulence reduces the rate at which the plumes transport sediment, and the magnitude of this effect depends on the relative contributions of individual settling and plume-driven transport. For larger particles, the plume contribution is modest compared with their already rapid individual settling. Suppressing the plumes therefore has only a small effect on the total settling velocity, which remains nearly independent of $\mathcal{A}$. For fine particles, the plume contribution accounts for most of the descent, so their settling rate is governed largely by the intensity of the plumes, which is itself sensitive to the turbulent state in which the instability develops. The descent of fine particles therefore cannot be predicted from their properties alone. It also depends on the turbulent state of the water into which they are released. This bears directly on ocean alkalinity enhancement, which favours fine grains because they dissolve more rapidly and therefore operates precisely in the size range whose transport is most sensitive to $\mathcal{A}$. Dissipation rates in the surface ocean span several orders of magnitude depending on wind, wave and buoyancy forcing, so the ocean turbulent state at the time of release may exert a first-order control on how long the material is retained in the surface layer, rather than being a detail that can simply be averaged over.

7.2. Caveats and future directions

Pre-existing turbulence influences the flow through the time at which the instability begins. This conclusion depends on how the turbulent field is constructed, and two features of that construction limit how broadly it can be generalized.

First, the phases in equation (2.13) are drawn at random and carry no physical information. Naturally evolved turbulence has phase relationships established by nonlinear interactions, which organize the flow into coherent vortices, whereas a randomly phased field reproduces the distribution of energy across horizontal scales but not the structures that would ordinarily accompany that distribution. The initial condition should therefore be interpreted as a field with an ocean-like spectrum rather than as a realization of ocean turbulence. Whether coherent vortices would change the outcome is an open question with a plausible argument in either direction. Organized structures might seed the instability more effectively at the scales it selects, or they might instead strain the emerging mode more persistently than a random field does. Initializing from a precursor simulation of decaying or forced turbulence would help resolve this question, as would using a field constructed from laboratory or observational measurements, provided that these measurements resolve the scales at which the instability operates.

Second, and more consequentially, the field is imposed once at $t = 0$ and is not sustained. Real mixed-layer turbulence could be continuously replenished by wind stress, wave breaking and surface buoyancy loss, whereas here the perturbation decays under viscosity and has no external energy source. This is what allows the initial field to act through the onset time. A sustained field would continue to strain the plumes after they form, and the

question our simulations cannot answer is whether that ongoing strain would further reduce transport or whether the plumes, once established, would remain robust to it. Our results argue that the plumes are robust in the absence of continued forcing, and continuously forced simulations are the natural next step for testing whether that robustness survives when the forcing is sustained.

The Schmidt numbers used here are below their oceanic values, as is unavoidable in direct simulation, and in particular the sediment and salt are assigned equal molecular diffusivities; this makes the comparison between the two turbulent diffusivities cleaner than it would otherwise be, but it removes a genuine physical asymmetry. The particles are monodisperse and non-interacting, whereas natural and engineered suspensions are polydisperse and, at sufficient concentrations, may flocculate, which alters the distribution of effective settling velocities (Balachandar & Eaton 2010; Rahmani *et al.* 2022). A Lagrangian treatment, in which individual particles are tracked and the drag, added-mass and pressure-gradient forces are solved for each particle (Chou *et al.* 2015; Chou & Shao 2016; Biegert *et al.* 2017), would relax these restrictions. Such an approach would allow a distribution of particle sizes to be represented within a single simulation. The present configuration contains no mean shear, which in the ocean surface layer would provide both an additional source of turbulence and a mechanism for suppressing the emerging structures (Yang *et al.* 2023, 2025).

Particle dynamics represent only half of the problem for ocean alkalinity enhancement, where the efficiency of a deployment is set by the competition between the time the particles reside in the mixed layer and the time required for them to dissolve. Alkalinity released below the mixed layer cannot equilibrate with the atmosphere until that water resurfaces (Renforth & Henderson 2017; Köhler *et al.* 2013; Eisaman *et al.* 2023), so material that descends before dissolving is largely ineffective. Coupling a shrinking-core dissolution model (Renforth & Kruger 2013) to the current Eulerian tracer framework would be the next step. The particle diameter would then shrink as dissolution proceeds and would in turn set the settling velocity, allowing the fraction of the added alkaline mineral that dissolves within the mixed layer to be determined.

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Author contributions C.-L. Liu conceived and designed the study, performed the simulations, developed and carried out the analysis, interpreted the results and wrote the manuscript. A. J. K. Yang suggested on the design of the simulation set, contributed to the interpretation of the results and revised the manuscript.

Appendix A. Sensitivity to the initial perturbation

The phases in equation (2.13) are drawn at random, so a given $\mathcal{A}$ specifies an ensemble of initial fields rather than a single one. Because the transition to turbulence in a stratified shear layer depends sensitively on the disturbance it grows from (Liu *et al.* 2022), some spread between realizations at fixed $\mathcal{A}$ is unavoidable. The question is therefore not whether such scatter exists, but whether it is small compared with the systematic dependence on $\mathcal{A}$ that we interpret.

We repeat the $D = 15 \mu\text{m}$, $\mathcal{A} = 8.3 \times 10^{-4}$ case with an independent set of random phases, drawn from the same uniform distribution, with the imposed spectrum and every other parameter unchanged. This case is chosen deliberately: it is the amplitude at which the frontal speed departs most conspicuously from the trend of its neighbours (figure 15*b*), and so the point at which a realization effect would be most likely to be mistaken for a physical response.

Figure 18 compares the two realizations (red symbols, black bar) with the amplitude sweep (blue). For each quantity we take the range spanned by the realizations at $\mathcal{A} = 8.3 \times 10^{-4}$ and express it as a fraction of the range spanned by the whole $D = 15 \mu\text{m}$ family across $\mathcal{A}$; this ratio is the percentage quoted in each panel, and it measures directly how much of the systematic response could be accounted for by the choice of phases alone. It is 2% for the onset time, 10% for the bulk effective settling velocity and 16% for the volume-averaged flux. The response to $\mathcal{A}$ is therefore between roughly six and fifty times larger than the scatter between realizations, and the conclusions of §3 and §6 do not rest on the particular phases drawn.

The frontal speed behaves differently. Its two realizations span 58% of the range covered by the entire family, so a single pair of realizations at one amplitude is more than half as widely separated as the extremes of the sweep, and the apparent maximum at $\mathcal{A} = 8.3 \times 10^{-4}$ lies within that spread. The non-monotonicity noted in §6.2 is therefore consistent with realization scatter and should not be read as a physical response. This supports the decision taken there to use W_{fr} only to corroborate the magnitude of the enhancement and its dependence on particle size, and to rely on equation (6.3) for the dependence on $\mathcal{A}$.

Two points follow for the interpretation of the results. Comparisons spanning the full range of $\mathcal{A}$ are robust: the systematic signal is several times the scatter, which is why the sweep covers nearly two decades rather than a narrow band. Differences between neighbouring amplitudes are not resolved by a single realization, and we accordingly draw no conclusion from them; every result reported in this paper rests on the trend across the full range, and each is separated from the realization scatter by a margin of at least a factor of six. Sensitivity to the initial disturbance is intrinsic to the transition (Liu *et al.* 2022) and cannot be eliminated, only outweighed, and the design adopted here ensures that it is.

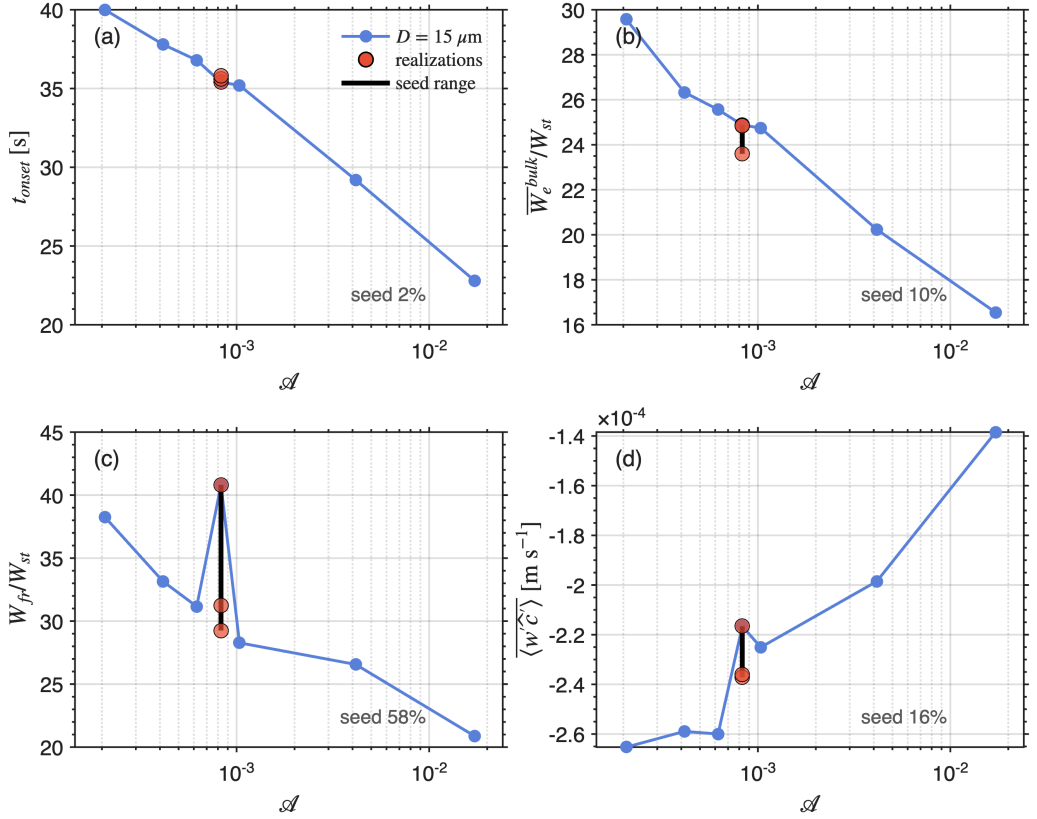


Figure 18. Sensitivity of the results to the random phases of the initial perturbation, for $D = 15 \mu\text{m}$. Blue curves show the amplitude sweep of table 2; red symbols mark three independent realizations at $\mathcal{A} = 8.3 \times 10^{-4}$, differing only in the phases drawn in equation (2.13), and the black bar spans their range. (a) Onset time, (b) bulk effective settling velocity, (c) frontal speed, and (d) volume-averaged sediment flux. The percentage quoted in each panel is the range spanned by the two realizations, expressed as a fraction of the range spanned by the entire $D = 15 \mu\text{m}$ family across $\mathcal{A}$; small values indicate that the systematic response to $\mathcal{A}$ is well separated from the scatter between realizations.

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