1 Flow splitting in numerical simulations of oceanic 2 dense-water outflows

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12 Abstract

Flow splitting occurs when part of a gravity current becomes neutrally buoyant and separates from the bottom-trapped plume as an interflow. This phenomenon has been previously observed in laboratory experiments, small-scale water bodies (e.g., lakes) and numerical studies of small-scale systems. Here, the potential for flow splitting in oceanic gravity currents is investigated using high-resolution ($\Delta x = \Delta z = 5$ m) two-dimensional numerical simulations of gravity flows into linearly stratified environments. The model is configured to solve the non-hydrostatic Boussinesq equations without rotation. A set of experiments is conducted by varying the initial buoyancy number $B_0 = Q_0 N^3 / g'^2$ (where Q_0 is the volume flux of the dense water flow per unit width, N is the ambient stratification and g' is the reduced gravity), the bottom slope (α) and the turbulent Prandtl number (Pr). Regardless of α or Pr , when $B_0 \leq 0.002$ the outflow always reaches the deep ocean forming an underflow. Similarly, when $B_0 \geq 0.13$ the outflow always equilibrates

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at intermediate depths, forming an interflow. However, when $B_0 \sim 0.016$, flow splitting always occurs when $Pr \geq 10$, while interflows always occur for $Pr = 1$. An important characteristic of simulations that result in flow splitting is the development of Holmboe-like interfacial instabilities and flow transition from a supercritical condition, where the Froude number (Fr) is greater than one, to a slower and more uniform subcritical condition ($Fr < 1$). This transition is associated with an internal hydraulic jump and consequent mixing enhancement. Although our experiments do not take into account three-dimensionality and rotation, which are likely to influence mixing and the transition between flow regimes, a comparison between our results and oceanic observations suggests that flow splitting may occur in dense-water outflows with weak ambient stratification, such as Antarctic outflows.

¹³ *Keywords:* intermediate and bottom water formation, internal hydraulic
¹⁴ jump, overflows, gravity currents

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1. Introduction

Oceanic outflows are large-scale density-driven currents generated either at high latitudes, where water density increases due to cooling and brine rejection from growing sea ice, or in subtropical marginal seas, where water density increases due to evaporation. These outflows have been widely studied (Legg et al., 2009; Price and Baringer, 1994), since they play an important role in establishing the large-scale circulation. In general, high latitude outflows, e.g. Denmark Strait outflow (Girton and Sanford, 2003), Faroe Bank Channel outflow (Mauritzen et al., 2005) and Antarctic outflows (Foldvik et al., 2004; Gordon et al., 2004; Muench et al., 2009), sink to the bottom of their respective basins, while outflows located closer to the equator, e.g., Mediterranean Sea outflow (Baringer and Price, 1997; Price et al., 1993) and Red Sea outflow (Matt and Johns, 2007; Peters et al., 2005), equilibrate at intermediate depths.

There is no clear observational evidence that major oceanic outflows, such as those discussed by Price and Baringer (1994) and Legg et al. (2009), can vertically split into two or more branches as a consequence of mixing and affect intermediate and deep layers simultaneously¹. Nevertheless, flow splitting in smaller scale density-driven currents has been observed in lakes (De Cesare et al., 2006), in a Mediterranean reservoir (Cortés et al., 2014a) and in the Arctic Ocean (Aagaard et al., 1985). Given the difficulties in

¹The Mediterranean outflow splits into two cores after the flow begin to descend into the Gulf of Cadiz, but this is a consequence of topographic steering (Baringer and Price, 1997). In the present work, focus is given to flow splitting induced by mixing rather than topographic effects.

56 measuring oceanic outflows, especially at high latitudes, and the fact that
 57 there are many processes acting in these flows concurrently (e.g., tides, hy-
 58 draulic control, thermobaric effect and geostrophic eddies), it is possible that
 59 flow splitting could also take place in large-scale oceanic outflows but has
 60 not been identified. If that is the case, conceptual models developed to close
 61 the heat budget in the global overturning circulation (Hughes and Griffiths,
 62 2006; Wåhlin and Cenedese, 2006; Wells and Wettlaufer, 2007) would have
 63 to be revisited, since they assume that large-scale high-latitude oceanic out-
 64 flows only ventilate the abyssal ocean. Throughout this paper, we follow the
 65 terminology used by limnologists and classify the different flow regimes as
 66 interflow, underflow or split flow.

67 Much of what is known about dense outflows comes from numerous labo-
 68 ratory experiments of bottom gravity currents (Baines, 2001, 2005; Cenedese
 69 et al., 2004; Cenedese and Adduce, 2008; Cortés et al., 2014b; Ellison and
 70 Turner, 1959; Lane-Serff and Baines, 1998; Mitsudera and Baines, 1992; Mon-
 71 aghan, 2007; Simpson, 1999; Wells and Wettlaufer, 2007; Wells and Nadara-
 72 jah, 2009). Ellison and Turner (1959) proposed a theory for the bulk prop-
 73 erties of a density gravity current flowing into homogeneous ambient water.
 74 They showed that the entrainment velocity of the plume (W_e) could be rep-
 75 resented as the product of the mean fluid velocity (U) and an entrainment
 76 parameter (E) that is an empirical function of the Richardson number (Ri)
 77 for the layer. Cenedese et al. (2004) added the effects of a rotating tank
 78 and found the development of three flow regimes (laminar, waves and ed-
 79 dies), where each of them has different mixing characteristics. In a follow-up
 80 study, Cenedese and Adduce (2008) showed that E also depends on the

Figure 1: An example of flow splitting observed in a non-rotating laboratory experiment where dense water flows down on a steep shelf slope ($\theta = 30^\circ$) and into a linearly stratified ambient water (adapted from Baines, 2008). The fluid has been dyed with fluorescence and is illuminated by a thin laser beam that scans a central vertical section.

81 Reynolds number (Re), and not just on Ri as previously proposed by Ellison
82 and Turner (1959).

83 For non-rotating gravity currents flowing down slopes and into linearly
84 stratified environments, Mitsudera and Baines (1992) showed the develop-
85 ment of a well-defined turbulent layer with up- and downslope streaming
86 formed above the bottom current. These experiments also revealed the pres-
87 ence of cusp-shaped waves near the top of the slope, which were attributed
88 to generation of Holmboe-like instabilities. Baines (2001) then proposed that
89 these instabilities are responsible for the wisps of detrained fluid observed in
90 laboratory experiments. These observations highlighted the different charac-
91 teristics of the flow compared with experiments with homogeneous ambient
92 stratification, and motivated additional laboratory experiments by Baines
93 (2005, 2008). In the latter, ranges of ambient stratification frequencies (N)
94 and slope angles (θ) were used and two main types of flow regimes emerged:
95 1) for sufficiently large θ , the flow was dominated by strong entrainment
96 balanced by buoyancy force; and 2) for small θ , the flow detrained and the
97 balance was mainly controlled by buoyancy force and bottom drag. Herein,
98 we use $\alpha = \tan(\theta)$ to represent the slope of θ .

99 An additional flow regime, characterized by the splitting of the gravity
100 current and the presence of both entrainment and detrainment, was also
101 observed by Baines (2005, 2008) (Fig. 1), but it was not explored in detail and
102 remains unexplained in continuously stratified environments. These different
103 flow regimes were investigated in the context of a non-dimensional parameter,
104 namely the initial buoyancy number (B_0):

$$B_0 = \frac{Q_0 N^3}{g_0'^2}, \quad (1)$$

105 where g_0' is the reduced gravity ($g_0' = g \frac{\Delta\rho}{\rho_0}$), Q_0 is the volume flux of the
 106 dense water per unit width at the top of the slope, g is the acceleration
 107 due to gravity, $\Delta\rho$ is the density difference between the outflow and the
 108 environmental fluid just above it, and ρ_0 is a reference density.

109 The flow splitting regime shown in Fig. 1 was observed in just two exper-
 110 iments with a very large angle ($\theta = 30^\circ$) and with intermediate values of B_0
 111 (1.7 and 4.0×10^{-3}). It is characterized by the splitting of the main plume
 112 into two branches: a denser and less diluted plume at the bottom, and a less
 113 dense and more diluted plume located at intermediate depths.

114 Density flows can also be studied in numerical models. Guo et al. (2014)
 115 conducted numerical simulations of non-rotating laboratory-scale gravity cur-
 116 rents descending a slope into a linearly stratified environment. Based on front
 117 positions, density and vertical velocity profiles, they showed that their sim-
 118 ulations were in good agreement with laboratory experiments (Baines, 2001,
 119 2005; Mitsudera and Baines, 1992). However, Guo et al. (2014) did not ex-
 120 plore the flow splitting regime since their simulations were only run for a
 121 short time and just one of their experiments showed the beginning of a split
 122 scenario with blobs of dense water detaching from the main current (see their
 123 Figs. 3c and d). While there have been some numerical studies investigating
 124 flow splitting in two-layer systems (Cortés et al., 2015; Wobus et al., 2013),
 125 this type of flow has never been modeled in detail in a linearly stratified
 126 environment.

127 Additional insights can be gained from studies based on non-rotating lab-

128 oratory experiments of dense water flowing down slopes into a two-layered
 129 ambient stratification (Cortés et al., 2014b; Monaghan, 2007; Wells and Wet-
 130 tlauffer, 2007). Three possible outcomes have been described for this setup
 131 (see Figs. 1a, b and c in Cortés et al., 2014b): 1) the flow separates from
 132 the bottom and forms an interflow; 2) the flow reaches the bottom of the
 133 domain and forms an underflow; or 3) the flow splits as it impinges into the
 134 single density step and forms two intrusions. The latter is the equivalent to
 135 the flow splitting in a linearly stratified environment shown in Fig. 1.

136 Motivated by observational evidence of flow splitting in a Mediterranean
 137 reservoir, Cortés et al. (2014b) conducted a series of laboratory experiments
 138 to investigate the splitting of a gravity current upon reaching a density in-
 139 terface. These authors proposed that flow splitting can be predicted based
 140 upon two non-dimensional parameters: the density Richardson number (Ri_ρ)
 141 and the densimetric Froude number (Fr). These numbers are defined as:

$$Ri_\rho = \frac{g'_{12} h_1}{B_f^{2/3}}, \quad (2)$$

142 and

$$Fr = \frac{U}{\sqrt{g' h}}, \quad (3)$$

143 where $g'_{12} = g(\rho_2 - \rho_1)/\rho_1$ is the reduced gravity of the ambient water density
 144 step; ρ_1 and ρ_2 are the densities of the upper and lower layer, respectively;
 145 h_1 is the thickness of the upper layer; $B_f = g(\rho_o - \rho_1)Q_0/(\rho_1 W)$ is the inflow
 146 buoyancy flux; ρ_o and W are the initial density and width of the outflow,
 147 respectively; U is the mean bottom current speed; and h is the outflow
 148 thickness. For Ri_ρ in the range $5 \leq Ri_\rho \leq 40$, flow splitting was more likely

149 to occur as the density interface of the gravity current became more diffusive
150 (i.e., when $Fr > 1$). In a follow-up study, Cortés et al. (2015) performed
151 numerical simulations to further explore the behavior of a gravity current
152 in a two-layer system. They found good agreement between simulations
153 and laboratory experiments, and confirmed the importance of the internal
154 properties of the gravity current in determining the different flow regimes.

155 The fate of dense outflows in nature is expected to depend on the shear
156 and density gradient at the plume’s interface (i.e., Fr) and the resulting
157 turbulence (i.e., Re) (e.g., Girton and Sanford, 2003). Thus, to apply the
158 results from laboratory experiments to geophysical flows these parameters
159 must be similar. In laboratory experiments, turbulence can only be generated
160 by using slope angles much larger than those observed in nature. Otherwise,
161 the density current cannot accelerate within the bounds of typical laboratory
162 tanks [$O(1)$ m] and the flow is not turbulent. Even when large slope angles
163 are used in laboratory studies, Re is still orders of magnitude smaller than
164 in geophysical flows (e.g., Baines, 2001). Therefore, it is not known whether
165 the results from laboratory experiments discussed above remain valid for
166 geophysical parameters and scales.

167 Since some realistic parameters cannot be reproduced in the laboratory,
168 numerical simulations have been used to complement our understanding of
169 the dynamics of bottom density currents from a geophysical point of view
170 (e.g., Ezer, 2005; Jackson et al., 2008; Legg et al., 2006, 2009; Özgökmen and
171 Chassignet, 2002; Özgökmen et al., 2004, 2006; Özgökmen and Fischer, 2008;
172 Wobus et al., 2013). The two-dimensional (2D) non-hydrostatic simulations
173 presented by Özgökmen and Chassignet (2002) captured qualitatively the

174 major features of bottom gravity currents observed in laboratory studies, e.g.,
 175 the evolution of the plume's head and its leading edge, and the development
 176 of Kelvin-Helmholtz instabilities. In a follow up study, Özgökmen et al.
 177 (2004) quantified the differences between 2D and three-dimensional (3D)
 178 simulations of bottom gravity currents. They found that the 3D results were
 179 in better agreement with laboratory experiments than the 2D counterpart.
 180 In particular, in the 2D experiments mixing due to entrainment was up to
 181 two times larger than in 3D. The authors attributed this discrepancy to the
 182 fact that 2D simulations cannot represent secondary instabilities that take
 183 place in 3D, resulting in a more rapid break down of the vortices. This
 184 difference was also observed in numerical simulations of gravity currents in
 185 the lock-exchange configuration (Härtel et al., 2000; Necker et al., 2002).

186 Özgökmen et al. (2006) used a combination of 2D and 3D numerical ex-
 187 periments to explore the effects of ambient stratification on the entrainment
 188 and properties of dense bottom currents that equilibrated at intermediate
 189 depths (i.e., forming interflows). Their simulations investigated the param-
 190 eter space of $0.125 \leq B_0 \leq 0.46$ and $3^\circ \leq \theta \leq 7^\circ$; it will be shown later
 191 that these values of B_0 are much larger than the range we find conducive to
 192 the development of flow splitting. Özgökmen et al. (2006) found that mix-
 193 ing by Kelvin-Helmholtz rolls did not penetrate all the way to the bottom
 194 of the density current and, therefore, the densest bottom water remained
 195 undiluted until the equilibration level was reached. They also demonstrated
 196 that entrainment and the properties of the product water depend only on the
 197 ambient stratification of the system, with no significant difference between
 198 2D and 3D experiments. However, as noted by the authors, the regime where

199 more vigorous mixing can occur remains to be investigated using a deeper
 200 domain, so that the mixing and consequent equilibration level with smaller
 201 B_0 can be captured. Wobus et al. (2013) used a hydrostatic model to per-
 202 form idealized three-dimensional simulations of a plume descending a conical
 203 slope into a rotating three-layer environment. By varying the inflow salinity
 204 and flow rate, they simulated flow regimes similar to those discussed above
 205 (interflow, underflow or split flow). Nonetheless, since their setup is based on
 206 a three layer system, it is not possible to classify their results in the context
 207 of B_0 .

208 Based on the preceding review, we have identified the need for additional
 209 simulations to develop our understanding of whether flow splitting can oc-
 210 cur in realistic oceanic outflows. Our new simulations cover the range of B_0
 211 for which we expect flow splitting, and use sufficiently high resolution (Δx
 212 $= \Delta z = 5$ m) and non-hydrostatic dynamics needed to resolve shear-driven
 213 instabilities, an important mechanism for mixing in bottom gravity currents
 214 (Simpson, 1999). Inspired by the large range (5 to 20) of averaged turbu-
 215 lent Prandtl number values ($Pr = \frac{\nu_h}{K_h} = \frac{\nu_v}{K_v}$, where ν_h/ν_v and K_h/K_v are
 216 the turbulent horizontal/vertical viscosity and diffusivity, respectively) re-
 217 ported by Cortés et al. (2015), we investigate the sensitivity of the solutions
 218 to the following Pr values: 1, 10 and 50. Our simulations reveal, for the
 219 first time, that the flow splitting regime may occur in non-rotating linearly
 220 stratified oceanic environments. A conceptual model for the development
 221 of flow splitting is proposed based on a particular class of shear instability,
 222 namely those in the transition zone between Kelvin-Helmholtz and Holmboe
 223 instabilities, and on the development of an internal hydraulic jump. Finally,

224 the conditions under which flow splitting might occur in oceanic outflows are
225 discussed.

226 We emphasize that numerous simplifications, summarized in section 4.4,
227 in our simulations preclude confirming that split flows actually occur in
228 oceanic outflows and make significant contributions to pathways for ventila-
229 tion of dense shelf water into the deep ocean. However, these results identify
230 a potentially important process that can be further investigated with more
231 sophisticated numerical simulations and micro- and fine-structure measure-
232 ment programs in candidate outflows.

233 **2. Methods**

234 *2.1. Numerical model and dynamical configuration*

235 The numerical simulations were conducted using the Massachusetts In-
236 stitute of Technology general circulation model (MITgcm) (Marshall et al.,
237 1997a,b). The model was configured to solve the non-hydrostatic Boussinesq
238 equations in the absence of rotation. MITgcm is a height-coordinate model
239 that employs partial step topography (Adcroft et al., 1997) and the config-
240 uration employed in this study includes a free surface. Given the available
241 computational resources, we only performed 2D simulations. Nevertheless,
242 the 2D approximation is a good background for future 3D rotating studies.
243 The relative inexpensiveness of 2D simulations allows for a more complete
244 parameter space investigation, so that the controlling physical processes can
245 be identified.

Figure 2: Schematic of the model domain.

246 2.2. Experimental setup and parameters

247 The model domain is depicted in Fig. 2. It includes a 5 km wide flat
 248 continental shelf that is connected to a flat deep ocean (which is 50 km wide)
 249 through a constant shelf slope (α) of width $L-5$ km, where L is the position
 250 where the slope intersects the ocean floor. The total length ($L+50$ km) varies
 251 between 70 and 115 km, depending on the slope used. We performed a set of
 252 experiments *a priori* to ensure that the flow characteristics described here are
 253 not influenced by the total length of the domain. Three different shelf slopes
 254 were used: steep ($\alpha = 0.1$), moderate ($\alpha = 0.05$) and shallow ($\alpha = 0.025$).
 255 This range covers the slopes of most continental shelves next to major oceanic
 256 outflows (Price and Baringer, 1994). The depth varies from 500 m at the shelf
 257 to 2000 m at the deep ocean. The geometry is consistent with the outer shelf
 258 and slope of the western Ross Sea, Antarctica (e.g., Padman et al., 2009),
 259 where shelf water depth is much larger than for most other global outflows.
 260 The dense water is prescribed in the first 2.5 km of the domain and below
 261 300 m depth (red hatched region in Fig. 2). This configuration leads, in our
 262 simulations, to an initial outflow thickness h_o at the top of the slope (i.e., x
 263 $= 5$ km) of $h_o \sim 100$ m. This is a reasonable representation of the initial
 264 thickness observed in the major oceanic outflows (Legg et al., 2009; Price
 265 and Baringer, 1994).

266 To avoid complications that may arise through processes such as thermo-
 267 baricity and cabbeling (McDougall, 1987), we ignore salinity and use a linear
 268 equation of state as a function of temperature (T) only: $\rho(T) = \rho_0(1 - \alpha_T T)$,
 269 where $\rho_0 = 1025 \text{ kg m}^{-3}$ is the background density and $\alpha_T = 5.3 \times 10^{-5} \text{ }^\circ\text{C}^{-1}$
 270 is the thermal expansion coefficient. A passive tracer (τ) with an initial con-

271 centration of 1 was imposed within the dense plume to help visualize the
 272 flow behavior. No-slip boundary conditions are applied to the momentum
 273 equations at the bottom, while no-flux conditions are applied to tracers at
 274 the free surface. The side boundaries are kept closed and tracers and mo-
 275 mentum are nudged towards the initial condition within regions that are 2.5
 276 and 5.0 km wide on the left and right boundaries, respectively (see Fig. 2). A
 277 relaxation timescale of 360 seconds is used for both tracers and momentum.
 278 This procedure prevents internal waves being reflected at the boundaries and
 279 also provides a source and sink for momentum and tracers.

280 Both horizontal (Δx) and vertical (Δz) resolutions are set to 5 m, since
 281 previous 2D non-hydrostatic simulations of bottom gravity currents designed
 282 to study geophysical parameters and scales (with $h_o \sim 100$ m) have shown
 283 that such grid spacing is small enough to resolve shear instabilities such as
 284 Kelvin-Helmholtz billows (Özgökmen and Chassignet, 2002). These billows
 285 pull lighter water down under denser water, and the mixing associated with
 286 convective readjustment results in water with intermediate density. This is
 287 an important process that takes place at the interface between the dense
 288 water and the ambient lighter water. We avoid spurious mixing that takes
 289 place near steplike topography of z-models by satisfying the condition $\Delta x <$
 290 $\Delta z/\alpha$ (Winton et al., 1998). Hence, our simulations are able to capture
 291 the homogenization of tracers induced by the mixing driven by frictionally
 292 generated shear. To check for convergence of our solutions as functions of
 293 resolution, selected experiments were run at $\Delta x = \Delta z = 2.5$ m and the results
 294 remained the same.

295 The initial ambient stratification (N) is constant throughout the water

296 column. Three different background density profiles are used: 5.0×10^{-4}
 297 s^{-1} (weakly stratified); $1.0 \times 10^{-3} \text{s}^{-1}$ (moderately stratified); and 2.5×10^{-3}
 298 s^{-1} (strongly stratified). In these cases, the density differences between the
 299 heaviest (i.e., ρ at $z = -2000 \text{ m}$) and lightest (i.e., ρ at $z = 0 \text{ m}$) water
 300 are 0.05, 0.2 and 1.3 kg m^{-3} , respectively. The initial density differences
 301 between the outflow and the ambient water just above it (i.e., ρ at $z = -$
 302 300 m) are set to $\Delta\rho = 0.2 \text{ kg m}^{-3}$ in the weakly and moderately stratified
 303 cases, and to 0.4 kg m^{-3} in the strongly stratified cases. Therefore, in the
 304 absence of mixing, the imposed outflow densities would result in underflows
 305 for the weakly and moderately stratified cases and interflows for the strongly
 306 stratified experiments. Since the initial development of the system resembles
 307 lock exchange flows, the initial volume flux of the dense water per unit width
 308 can be computed as $Q_0 \sim \sqrt{g'h_0}h_0$ (e.g., Özgökmen et al., 2004). With $h_0 \sim$
 309 100 m and using the values of N and $\Delta\rho$ listed above, we vary B_0 (as defined
 310 in Eq. 1) by two orders of magnitude: 0.002, 0.016 and 0.13.

311 Laplacian friction was employed in the momentum equations. The ver-
 312 tical component of viscosity was set to $\nu_v = 10^{-4} \text{ m}^2 \text{s}^{-1}$, while the hori-
 313 zontal component (ν_h) was $7.5 \text{ m}^2 \text{s}^{-1}$ in the strongly stratified cases, and
 314 $5.0 \text{ m}^2 \text{s}^{-1}$ otherwise. These values were chosen such that the grid Reynolds
 315 number ($\text{Re}_g = U_{max}\Delta x/\nu_h$, where U_{max} is the maximum plume velocity) was
 316 always less than 1.5. This condition is effective in eliminating grid-scale noise
 317 (Griffies et al., 2000; Ilicak et al., 2012), and is needed because momentum
 318 is advected using a second-order scheme that does not include an equivalent
 319 constraint. The horizontal viscosity values employed here are much larger
 320 than laboratory and realistic oceanic conditions, but they are required for

numerical reasons. Our simulations are, therefore, at lower Reynolds number than laboratory and true oceanic flows.

Unlike momentum, tracers are advected using a third-order direct space-time flux limited scheme, which eliminates grid-scale noise by introducing numerical diffusivity where needed. Since the problem investigated here includes shocks and propagating fronts, the use of a flux limited scheme is essential to avoid unphysical behavior in terms of mixing (e.g., Ilicak et al., 2012; Marques and Özgökmen, 2014). Although this scheme does not require explicit tracer diffusivities, we set the horizontal (K_h) and vertical (K_v) diffusivities so that each experiment is run with the turbulent Prandtl number ($Pr = \frac{\nu_h}{K_h} = \frac{\nu_v}{K_v}$) set to 1, 10 and 50. In stratified flows, momentum tends to diffuse faster than tracers, leading to $Pr > 1$. The reasoning behind varying Pr is to investigate how the flow behaves when the internal properties of the gravity current are changed. The values employed here are similar to the average Pr values (5 to 20) obtained in the numerical study of Cortés et al. (2015). Based on a set of experiments performed *a priori*, we verified that $Pr = 50$ is the approximate maximum value that can be achieved in our experiments (i.e., when K_h is approximately equal to the implicit numerical diffusion). Ideally, Pr should be represented in terms of the gradient Richardson number ($Ri(z) = \frac{N^2}{S^2}$, where $S = \frac{\partial u}{\partial z}$ is the vertical shear) using physically based formulations (e.g., Strang and Fernando, 2001; Elliott and Venayagamoorthy, 2011). However, none of these formulations has been implemented in the model employed here and, therefore, we choose to explicitly set the values of Pr .

Using the above mentioned values of B_0 , Pr and α , a total of 27 2D

Table 1: Summary of parameters used in the numerical experiments.

Domain total length	70, 90 and 115 km
Grid resolution	$\Delta x = \Delta z = 5$ m
Maximum ocean depth	$H_{max} = 2000$ m
Shelf depth	$H_{shelf} = 500$ m
Linear equation of state	$\rho(T) = 1025(1 - \alpha_T T)$
Thermal expansion coefficient	$\alpha_T = 5.3 \times 10^{-5} \text{ } ^\circ\text{C}^{-1}$
Shelf slope	$\alpha = 0.025, 0.05$ and 0.1
Outflow thickness at $x = 5$ km	$h_0 \sim 100$ m
Ambient buoyancy frequency	$N = 5.0 \times 10^{-4}, 10^{-3}$ and $2.5 \times 10^{-3} \text{ s}^{-1}$
Initial $\rho_{outflow} - \rho_{ambient}$	$\Delta\rho = 0.2$ and 0.4 kg m^{-3}
Buoyancy number at $x = 5$ km	$B_0 = (2, 16 \text{ and } 130) \times 10^{-3}$
Vertical and horizontal viscosities	$\nu_v = 10^{-4} \text{ m}^2 \text{ s}^{-1}$; $\nu_h = 5$ and $7.5 \text{ m}^2 \text{ s}^{-1}$
Turbulent Prandtl number	$Pr = 1, 10$ and 50
Time step	$\Delta t = 0.25$ s

numerical experiments were conducted. All simulations were started from rest and the model was stopped once it became possible to identify whether the flow has formed an interflow, underflow or split flow. In all experiments, the time step (Δt) was set to 0.25 s. This value ensured that both the CFL condition and the stability parameter for the Laplacian friction were always satisfied. The model parameters are presented in Table 1.

3. Results

3.1. Initial flow development

The initial flow development of the system is evaluated using snapshots of passive tracer (τ) concentration and density contours at different instants

Figure 3: Time evolution of passive tracer distribution (τ , color) and density contours (kg m^{-3} , black lines plotted every 0.02 kg m^{-3}) for experiments with $\alpha = 0.1$, $B_0 = 0.002$ and $Pr = 1, 10$ and 50 (left, center and right panels, respectively). The initial evolution of the plume's head are shown at the top panels, intermediate states are shown in the middle panels and the conditions at the end of the simulations are shown in the bottom panels. The underflow regime is always generated, regardless of Pr . The blue boxes in the top and middle panels denote the boundaries for the closeups shown in the plots. Only a portion of the computational domain is displayed. Corresponding animations are available in the auxiliary materials (ms01, ms02 and ms03 refer to experiments with $Pr = 1, 10$ and 50 , respectively).

356 after each simulation was started. In the next two sections we focus on the
 357 nine experiments with $\alpha = 0.1$ (steepest model slope). With this slope, we
 358 varied the initial buoyancy number, $B_0 = 0.002$ (Fig. 3), 0.016 (Fig. 4) and
 359 0.13 (Fig. 5); and the turbulent Prandtl number, $Pr = 1, 10$ and 50 (on each
 360 figure: left, center and right panels, respectively).

361 The initial development of the system is similar to that of the so-called
 362 lock-exchange flow (e.g., Simpson, 1999), in which the denser fluid propagates
 363 downslope while the lighter fluid flows in the opposite direction along the top
 364 boundary of the domain. All cases show the development of a characteristic
 365 head as the leading edge of the outflow propagates downslope (panels a-c
 366 on each of Figs. 3 to 5). The initial growth of the head was observed in
 367 all experiments, as less dense ambient water was entrained into the leading
 368 edge of the plume. As Pr increases (i.e., when the diffusivities are reduced),
 369 the plume becomes less diluted and the vertical gradient of τ at the tail of
 370 the current (i.e., the region behind the head) becomes sharper (panels a-c on
 371 each of Figs. 3 to 5).

372 Once the head of the current has passed $x \sim 15$ km, density interfaces at
 373 the tail of the current begin to roll up (panels d-f on each of Figs. 3 to 5).
 374 These counter-clockwise vortices are indicative of shear instability, in which
 375 unstable waves entrain less dense fluid into the dense current. Both the
 376 horizontal and vertical dimensions of these features vary between 50 and 200
 377 m and, therefore, are resolved by the grid resolution employed here ($\Delta x = \Delta z$
 378 $= 5$ m). There is a general tendency for these vortices to coalesce and form
 379 larger structures via pairing, which is a well-documented behavior of 2D shear
 380 instabilities (e.g., Özgökmen and Chassignet, 2002). These vortices would be

Figure 4: Same as Fig. 3, but for $B_0 = 0.016$. The interflow regime occurs when $Pr = 1$ (left column), while flow splitting occurs when $Pr \geq 10$ (middle and right columns). Corresponding animations are available in the auxiliary materials (ms04, ms05 and ms06 refer to experiments with $Pr = 1$, 10 and 50, respectively).

381 significantly smaller if the effects of three-dimensionality were included in the
 382 simulations (Härtel et al., 2000; Özgökmen et al., 2004), potentially affecting
 383 the mixing of the flow. At this point, the initial development of clockwise
 384 vortices due to frictionally generated shear is also evident in the bottom
 385 boundary layer (red region in the inset of panels d-f on each of Figs. 3 to 5;
 386 see also correspondent animations in the Auxiliary Material). These features
 387 act as sinks of kinetic energy and, unlike the upper shear instability at the
 388 interface, tend to homogenize the water properties next to the bottom (i.e.,
 389 they have less effect on mixing of tracers). Further analyses of the different
 390 types of shear instabilities taking place in these simulations are presented in
 391 section 3.4.

392 3.2. *Flow regimes*

393 In the simulations where $B_0 = 0.002$ (Fig. 3), the initial head remains
 394 attached to the bottom current until it reaches the deep ocean (Figs. 3 d
 395 to f). Snapshots taken at the end of the simulations clearly show the gen-
 396 eration of an underflow, regardless of the value of Pr (Figs. 3 g to i). The
 397 only differences are that the plume becomes less diluted and the size of the
 398 instabilities gets smaller as Pr increases.

399 When $B_0 = 0.016$ (Fig. 4), the head merges with the trailing current
 400 before the leading edge reaches the end of the slope (Figs. 4 d to f). Soon
 401 after, the behavior of the current differs among the simulations. For $Pr =$
 402 1, the entire plume detaches from the bottom forming an interflow (Fig. 4g).
 403 However, when $Pr = 10$ or 50, a portion of the upper part of the current
 404 detaches from the slope and flows out into the interior, while the lower part
 405 of the current remains flowing towards the bottom of the domain (Figs. 4 h

Figure 5: Same as Fig. 3, but for $B_0 = 0.13$. The interflow regime is always generated, regardless of Pr . Density contours (black lines) are plotted every 0.1 kg m^{-3} . Corresponding animations are available in the auxiliary materials (ms07, ms08 and ms09 refer to experiments with $Pr = 1$, 10 and 50, respectively).

406 and i, respectively). In both cases, a clear discontinuity (or shock) separating
 407 two different mixing regions occurs around $x \sim 18$ km. The onshore region is
 408 characterized by a “pool” of homogeneous fluid (light blue shading in Figs. 4h
 409 and i) sitting on the top of the bottom current. This fluid has been strongly
 410 diluted by shear instabilities and subsequent mixing. The offshore region is
 411 characterized by thickening of the plume, with water classes ranging between
 412 the less dense intermediate plume to the denser and less diluted plume next
 413 to the bottom. This sudden increase in the plume thickness suggests the
 414 development of an internal hydraulic jump, and we explore this further in
 415 section 3.4.2. From Figs. 4 h and i, it is also evident that, as Pr increases,
 416 the lower part of the current is less diluted and, therefore, it penetrates
 417 to a greater depth (e.g., the distance between the upper and lower plumes
 418 increases). Nonetheless, in both cases the structure of the system resembles
 419 the flow splitting regime observed by Baines (2008) (see Fig. 1).

420 Finally, when $B_0 = 0.13$ (Fig. 5), for all values of Pr the head of the plume
 421 detaches from the bottom before the outflow reaches the end of the slope
 422 (Fig. 5 d to f). This plume then propagates towards the interior, resulting in
 423 an interflow (Figs. 5 g to i). As in the cases described above, as Pr increases,
 424 the gravity current gets less diluted and the spatial scale of features become
 425 smaller. There is a clear distinction between the lower and upper part of the
 426 flow when $Pr \geq 10$, where the fluid is significantly more diluted in the latter
 427 (Figs. 5 h and i). This distinction becomes more pronounced as Pr increases.
 428 However, since N in these experiments is relatively large, this condition is
 429 not enough to result in flow splitting.

430 3.3. Sensitivity to shelf slope

431 Laboratory experiments by Baines (2001, 2005) show that the bottom
432 slope plays a key role in determining the flow regime of downslope flows into
433 stratified environments. Therefore, it is useful to verify if the same reasoning
434 applies to the parameter space employed in our simulations. To this end,
435 we present additional simulations using $\alpha = 0.05$ and 0.025 . The remaining
436 parameters in these runs, namely B_0 and Pr , are identical to those described
437 in section 3.1. Hence, a total of eighteen additional experiments are described
438 here. From snapshots of τ taken at the end of these simulations, a qualitative
439 description of the different flow regimes is presented below.

440 As in the steep slope cases, when $B_0 = 0.002$ and $\alpha = 0.05$ (Figs. 6 a to c)
441 or $\alpha = 0.025$ (Figs. 7 a and c), the underflow regime is generated regardless
442 of Pr . Again, the only differences between these cases is that the plume
443 gets less diluted and smaller length-scale features become apparent as Pr
444 increases.

445 When $B_0 = 0.016$, $Pr = 1$ and $\alpha = 0.05$ (Figs. 6d) or $\alpha = 0.025$ (Figs. 7d),
446 the current detaches from the bottom and forms an interflow. For $Pr \geq 10$,
447 flow splitting also occurs under a moderate slope $\alpha = 0.05$ (Figs. 6 e and
448 f) and under a shallow slope $\alpha = 0.025$ (Figs. 7 e and f). In comparison to
449 the steep slope cases (Figs. 4h and i), as the slope decreases the plume needs
450 to travel further to result in flow splitting. A thin “tongue” of intermediate
451 density appears between the upper and lower parts of the plume (Fig. 6e).
452 This feature is not present when $\alpha = 0.1$.

453 Finally, when $B_0 = 0.13$, all simulations lead to the interflow regime,
454 regardless of α and Pr .

455 In summary, for given B_0 and Pr values, the flow regimes obtained using
 456 intermediate ($\alpha = 0.05$) and shallow ($\alpha = 0.025$) slopes are identical to those
 457 obtained using a steep slope ($\alpha = 0.1$). In the following sections, we focus
 458 on simulations conducted with the steep slope ($\alpha = 0.1$) since less disk space
 459 was required to save high frequency output, used in the following analysis,
 460 during the initial development of the flow in these runs.

461 3.4. *Internal properties of the outflow*

462 3.4.1. *Buoyancy and velocity profiles*

463 We now examine how buoyancy and cross-slope velocity profiles vary
 464 among our simulations. We define buoyancy as $\bar{b}(z) = -g(\bar{\rho}(z) - \rho_0)/\rho_0$, and
 465 cross-slope velocity, $\bar{u}(z)$, profiles, with over-lines indicating averages over
 466 space and time. These profiles were computed by first extracting $\rho(z, x, t)$
 467 and $u(z, x, t)$ for $5.5 \leq x \leq 6.0$ km (top of slope) then spatially averaging
 468 the data based on the distance of each level with respect to the bottom.
 469 Next, the data were averaged over time for $15 \text{ h} \leq t \leq 20 \text{ h}$ with an output
 470 period of 120 seconds. Since shear instabilities become more pronounced at
 471 later times and downstream of this region, this procedure provides a robust
 472 representation of the profiles, with relatively small standard deviations from
 473 the mean values.

474 As Pr increases, the buoyancy gradient ($\partial\bar{b}/\partial z$) becomes sharper while
 475 the shear ($\partial\bar{u}/\partial z$) does not change significantly throughout most parts of the
 476 water column (Fig. 8). When $Pr = 1$, the upper part of the profile (height
 477 ~ 80 m above bottom) displays a smoother shape than the corresponding
 478 profiles with $Pr = 10$ and 50. For a fixed Pr value, both $\bar{u}(z)$ and $\bar{b}(z)$ are
 479 overall very similar among the experiments. The only difference is that $\Delta\bar{u}$

Figure 6: Snapshots of passive tracer distribution (τ , color) and density contours (kg m^{-3} , black lines) taken at the end of the simulations for experiments with $\alpha = 0.05$. Flow splitting occurs when $B_0 = 0.016$ and $Pr \geq 10$ (panels e and f). The following parameters are varied: (from left to right) $Pr = 1, 10$ and 50 ; (from top to bottom) $B_0 = 0.002, 0.016$ and 0.13 . Density contours are plotted every 0.02 kg m^{-3} in the top and middle panels, and every 0.1 kg m^{-3} in the bottom panels.

Figure 7: Same as Fig. 6, but for experiments with $\alpha = 0.025$. Flow splitting occurs when $B_0 = 0.016$ and $Pr \geq 10$ (panels e and f).

Figure 8: Vertical profiles of cross-slope velocity ($\bar{u}$, black lines) and buoyancy ($\bar{b}$, red lines) for simulations with $B_0 = 0.002, 0.016$ and 0.13 (from top to bottom, respectively), and $Pr = 1, 10$ and 50 (from left to right, respectively). These plots highlight how the vertical gradient in buoyancy ($\partial\bar{b}/\partial z$) increases (i.e., the buoyancy profile gets sharper) as Pr increases. Height is shown as meters above bottom (mab). Over-lines indicate averages over space and time, see text for further details. Black and red stars represent the profiles fitted using Eqs. 4 and 5, respectively.

480 and $\Delta\bar{b}$ (defined in Fig. 9) are slightly larger in cases where $B_0 = 0.13$.

481 To quantify the differences among profiles, the following idealized curves
 482 for $\bar{u}(z)$ and $\bar{b}(z)$, respectively, are fitted to the model data using the non-
 483 linear least squares algorithm available in the open source Python library
 484 SciPy²:

$$\tilde{u} = -\frac{\Delta u}{2} \tanh\left(\frac{2}{h_u} z\right); \quad (4)$$

$$\tilde{b} = \frac{\Delta b}{2} \tanh\left(\frac{2}{h_b} z\right), \quad (5)$$

485 where h_u is the thickness of the shear layer, z is the vertical coordinate
 486 centered at the level of maximum shear, Δu and Δb are the velocity and
 487 buoyancy differences, respectively (see Fig. 9). Profiles fitted to the model
 488 data using Eqs. 4 and 5 are also shown in Fig. 8 (black and red stars, respec-
 489 tively). The idealized curves are not meant to capture boundary layer shears
 490 (which is of less importance here) but to represent the environment in which
 491 interfacial instabilities develop. A few important dimensionless parameters
 492 for the evolution of the flow can now be defined: the scale ratio, $R = h_u/h_b$,
 493 where h_b is thickness of the stratified layer (see Fig. 9); the bulk Richardson
 494 number, $J = \Delta b h_u / \Delta u^2$; and the central Richardson number, $Ri_0 = JR$.
 495 The central Richardson number quantifies the relative importance of strat-
 496 ification and shear and it is equivalent to the gradient Richardson number,
 497 $Ri(z) = \frac{\partial b / \partial z}{(\partial u / \partial z)^2}$ evaluated at the center of the shear layer. An additional
 498 parameter that is usually considered in studies of stratified shear flows (e.g.,

²http://docs.scipy.org/doc/scipy-0.16.1/reference/generated/scipy.optimize.curve_fit.html

499 Smyth et al., 2007) is the asymmetry factor (a), which quantifies the offset
 500 between the centers of the shear and stratified layers. However, we neglect
 501 this parameter since the offset in all profiles investigated here is very small
 502 (maximum of ± 5 m, or one model grid cell).

Figure 9: Schematic representation of (a) velocity and (b) buoyancy profiles constructed using Eqs. 4 and 5, respectively. Also shown are the important dimensionless parameters, R and Ri_0 , used to characterize the evolution of the flow (see text for their definition).

503 For all profiles shown in Fig. 8, R increases as Pr increases (Fig. 10a).
 504 The most significant change occurs when Pr varies between 1 and 10, where
 505 R approximately doubles (i.e., $h_u \sim 2 \times h_b$). As Pr changes from 10 to 50,
 506 the corresponding increase in R is relatively small ($\sim 20\%$) and at this point
 507 h_b is approaching the smallest value that we can resolve (~ 15 m, or $3 \times \Delta z$)
 508 given the vertical resolution employed here ($\Delta z = 5$ m). The thickness of the
 509 shear layer (h_u) decreases as Pr increases (Fig. 10b), while Ri_0 gets larger as

510 Pr increases (Fig. 10b). Again, the most significant changes in h_u and Ri_0
 511 occur in the $Pr = 1$ to $Pr = 10$ transition. In the absence of instabilities and
 512 given a time scale λ , both thicknesses will grow by diffusion according to h_u
 513 $\sim (\nu_v \lambda)^{1/2}$ and $h_b \sim (K_v \lambda)^{1/2}$. However, after 20 hours and given the small
 514 values of ν_v and K_v employed here (see Table 1), h_u and h_b cannot grow to
 515 the observed values just because of diffusion. Although instabilities in the
 516 upper part of the slope were not pronounced during the period when the
 517 profiles were averaged (between 15 and 20 hours), we speculate that h_u and
 518 h_b develop due to the horizontal advection of buoyancy and momentum in
 519 a sloped bottom. This is a consequence of the vertical coordinate employed
 520 here and the results might differ under different vertical coordinates, such as
 521 terrain-following or isopycnals.

522 The necessary condition for the development of instabilities based on
 523 the Miles-Howard criterion (Howard, 1961; Miles, 1961), namely $Ri(z) <$
 524 0.25 , is satisfied in all cases (Fig. 10c); recall that Ri_0 is equivalent to $Ri(z)$
 525 evaluated at the center of the shear layer. Although this condition is not
 526 sufficient for the development of instabilities, the results presented in the
 527 previous sections clearly show that instabilities occur in all simulations. It
 528 is not clear, however, whether these instabilities are all of the same type.
 529 There is a continuous spectrum of modes that range from Kelvin-Helmholtz
 530 to Holmboe instabilities and beyond, and it is not trivial to identify them
 531 (this is also beyond the scope of this paper). However, the fact that R changes
 532 significantly as Pr increases from 1 to 10 suggests that different instabilities
 533 might be occurring (see, e.g., Smyth and Winters, 2003; Hogg and Ivey, 2003)
 534 and we will discuss this further in Section 4.1.

34
Figure 10: Prandtl number versus (a) scale ratio, R ; (b) shear thickness, h_u [m] ; and (c) central Richardson number, Ri_0 . Results from simulations with different B_0 values are shown.

535 3.4.2. Internal hydraulic jump

536 The presence of internal hydraulic jumps is often associated with discon-
537 tinuities in the density and velocity fields, where the flow transitions from
538 supercritical to subcritical conditions (e.g., Holland et al., 2002; Legg and
539 Klymak, 2008; Long, 1953, 1954). Supercritical flows are characterized by
540 $Fr \geq 1$, i.e., the mean bottom current speed U exceeds the long-wave phase
541 speed $c_p = \sqrt{g'h}$, while the flow is subcritical when $Fr < 1$. It is therefore
542 of interest to investigate whether and why internal hydraulic jumps occur
543 in the different flow regimes simulated here. Following Wobus et al. (2013),
544 the thickness of the plume, $h(x, t)$, is defined as the height above the bot-
545 tom where the passive tracer concentration τ drops below 50% of its value
546 at the seabed. This selection is somewhat arbitrary and $h(x, t)$ is sensitive
547 to the chosen threshold value. Nonetheless, the outcomes shown below are
548 still valid when the upper threshold of τ was varied within 35% to 65% of
549 the bottom value. In the following analysis, we focus on experiments where
550 $\alpha = 0.1$ and $Pr = 50$ since the internal hydraulic jump is more pronounced
551 under these conditions. Just the region within the shelf slope is evaluated
552 and the values are averaged in time over the last 10 hours of each simulation
553 (denoted by angle brackets, $\langle \rangle$).

554 When an underflow is generated ($B_0 = 0.002$), $\langle h \rangle$ is almost constant
555 throughout the slope (Fig. 11a). Similarly, when the plume results in an
556 interflow ($B_0 = 0.13$), $\langle h \rangle$ remains constant until the current detaches from
557 the slope (Fig. 11a). However, when flow splitting occurs ($B_0 = 0.016$),
558 there is a strong discontinuity in $\langle h \rangle$ at $x \sim 16$ km (Fig. 11a). Indeed, the
559 time-averaged Froude numbers, $\langle Fr \rangle$, in the flow splitting regime show a

Figure 11: Time-averaged plume thickness, $\langle h_0 \rangle$ (a) and Froude number, $\langle Fr \rangle$ (b) as a function of cross-slope distance for selected experiments where $\alpha = 0.1$, $Pr = 50$ and $B_0 = 0.002$ (underflow), 0.016 (split flow) and 0.13 (interflow).

560 transition from supercritical to subcritical, consistent with the discontinuity
 561 in $\langle h \rangle$ (Fig. 11b). In contrast, $\langle Fr \rangle$ remains supercritical in the interflow and
 562 underflow regimes (Fig. 11b; noticed that the purple line drops off to zero as
 563 the plume detaches from the slope, as shown in Fig. 11a).

564 To further investigate what is causing the internal hydraulic jump in the
 565 flow splitting regime, we compute the density excess of the system as follows:

$$\rho^*(x, z, t) = \rho(x, z, t) - \rho(x, z, t = 0), \quad (6)$$

566 where $\rho(x, z, t = 0)$ is the initial density field. Time-averaged density excess
 567 $\langle \rho^* \rangle$ for a simulation that results in flow splitting ($B_0 = 0.016$) clearly shows a
 568 wide region with densities smaller than those at the start of the run (Fig. 12).
 569 This region coincides with the pool of homogeneous τ described earlier and
 570 shown in Fig. 4i. One possible explanation for the formation of this buoyant
 571 region, where $\langle \rho^* \rangle$ is negative, is the development of Holmboe-like instabilities
 572 upstream of this region. These instabilities engulf parcels from the upper part
 573 of the gravity current leading to detrainment. The detrained parcels of fluid
 574 are initially advected downstream due to inertia, leading to further mixing
 575 with ambient water (i.e., entrainment). The result is that the detrained
 576 water parcels are brought to depths greater than those where they would be
 577 neutrally buoyant.

578 This previous description is supported by the time evolution of 1444 par-
 579 ticles advected offline using a fourth-order Runge-Kutta scheme (Fig. 13).
 580 The particles were initially seeded between $14 \text{ km} \leq x \leq 20 \text{ km}$ and -2000
 581 $\text{m} \leq z \leq -800 \text{ m}$, every 150 m in the horizontal direction and 25 m in the
 582 vertical direction (Fig. 13a, see also corresponding animation available in the

Figure 12: Time-averaged density excess $\langle \rho^* \rangle$ as a function of depth and cross-slope distance for an experiment where flow splitting occurs ($\alpha = 0.1$, $Pr = 50$ and $B_0 = 0.016$). Blue contours show regions where water is less dense than the initial ambient water at the same depth; i.e., it indicates where water parcels are brought to depths greater than those where they would become neutrally buoyant.

583 auxiliary materials). The velocity field was updated every 120 seconds and
 584 the particles were advected between $15 \text{ h} \leq t \leq 45 \text{ h}$. The movement of
 585 particles reveals a slow counter-clockwise circulation on the top of the den-
 586 sity current, between approximately $10 \text{ km} \leq x \leq 18 \text{ km}$. This is also the
 587 region where $\langle \rho^* \rangle$ is negative (Fig. 12), which reinforces the previous expla-
 588 nation that detrained parcels of fluid are advected downstream of their level
 589 of neutral buoyancy, and then rise due to buoyancy.

590 4. Discussion

591 4.1. Conceptual model for flow splitting in linearly stratified environments

592 Previous sections have elucidated the different flow regimes that may
 593 occur in numerical simulations of oceanic dense-water outflows with linearly
 594 stratified environments. However, a conceptual model for the development
 595 of the flow splitting regime under such conditions is still missing. This model
 596 is now presented, followed by a discussion on how it compares with processes
 597 previously suggested in the literature. The description that follows applies
 598 to all experiments where $B_0 = 0.016$ and $Pr \geq 10$, regardless of α .

599 A summary of the processes occurring in the flow splitting regime is
 600 depicted in Fig. 14. The outflow is initially supercritical ($Fr > 1$), with
 601 shear-generated instabilities developing at the top boundary of the current
 602 (dark blue circles in Fig. 14). Additional shear-induced instabilities also
 603 occur in the bottom boundary layer (yellow circles in Fig. 14), but these act
 604 as sinks of kinetic energy and just homogenize the dense plume, i.e., they
 605 have less effect on mixing than the instabilities in the upper interface. The
 606 upper instabilities are characterized by buoyancy interfaces that are thinner

Figure 13: Time evolution of 1444 particles advected between $15 \text{ h} \leq t \leq 45 \text{ h}$ using the velocity field from an experiment where flow splitting occurs ($\alpha = 0.1$, $Pr = 50$ and $B_0 = 0.016$). The particles were seeded at time = 15 hours and their time evolution reveals a slow counter-clockwise circulation on the top of the density current, between approximately $10 \text{ km} \leq x \leq 18 \text{ km}$. Snapshots of passive tracer distribution (τ , color) are also plotted in the background. Corresponding animation is available in the auxiliary materials (ms10).

Figure 14: Schematic depiction of the processes occurring in the flow splitting regime.

607 than the shear interfaces (i.e., $R > 2$), which is suggestive of Holmboe-like
 608 instabilities (e.g., Smyth and Winters, 2003; Hogg and Ivey, 2003). They lead
 609 to detrainment of small blobs of fluid that are then advected downstream, as
 610 a consequence of inertia, to form a “pool” of stagnant buoyant fluid sitting
 611 on the top of the main current (Fig. 13; also highlighted by the light blue
 612 region in Fig. 14). This buoyant water is brought to depths below its level
 613 of neutral buoyancy relative to the environment (Fig. 12). Since there is a
 614 constant injection of detrained fluid upstream, the presence of a region of
 615 buoyant fluid results in a slow counter-clockwise circulation on the top of
 616 the density current (Fig. 13; also highlighted in Fig. 14). A small fraction of
 617 fluid then propagates towards the deep ocean away from the slope, forming
 618 the intermediate plume as the fluid becomes neutrally buoyant, while the
 619 rest of the fluid recirculates back towards the upper part of the slope. As
 620 the main density current propagates underneath the region of stagnant fluid,
 621 it gets thicker but its velocity is approximately constant. Both interfacial
 622 and near-bottom shear instabilities also occur at this state. The flow then
 623 transitions from a supercritical condition ($Fr > 1$) to a slower and more
 624 uniform subcritical condition ($Fr < 1$). This transition results in an internal
 625 hydraulic jump, where kinetic energy is converted into potential energy as a
 626 consequence of strong mixing (red region in Fig. 14). Additional water masses
 627 are then generated, filling the gap between the upper and lower plumes. The
 628 fluid in the bottom boundary layer undergoes little mixing with ambient
 629 water, so that the lower abyssal plume is almost undiluted.

630 There are many similarities between the mechanism described above and
 631 the scenario described by Baines (2008) based on laboratory experiments.

632 Baines (2008) suggested that detrainment in the upper part of the current
 633 could lead to the development of the intermediate plume observed in Fig. 1.
 634 Detrainment has been previously described as a consequence of Holmboe-like
 635 instabilities generated due to the mutual interaction between a gravity wave
 636 on the interface and a vorticity wave traveling above the density current
 637 (Baines, 2001). As discussed above, detrainment of small blobs of fluid in
 638 combination with the characteristics of the internal properties of the flow
 639 also suggest the presence of Holmboe-like instabilities in our simulations.
 640 However, unlike previous studies (e.g., Baines, 2001), the offset between the
 641 interfacial wave and the line of maximum vorticity is insignificant in our
 642 simulations (maximum of 5 m or one grid cell). Rather, the instabilities
 643 in our simulations are triggered as a consequence of the buoyancy interface
 644 being sharper than the shear interface (i.e., $R > 2$). Consistent with our
 645 numerical results, estimates of the Richardson number in the two laboratory
 646 experiments that resulted in flow splitting were more than two times larger
 647 than those experiments where splitting did not occur (see Table 1 in Baines
 648 (2005)). Based on a “conspicuous upward bubbling” (Fig. 1), Baines (2008)
 649 suggested that the fluid supplying the intermediate plume arose from the
 650 vicinity of the bottom plume. This indicates that portions of the plume were
 651 advected to depths greater than those where they would be neutrally buoyant
 652 and, therefore, the fluid rose due to buoyancy. This is also consistent with
 653 our simulations where flow splitting occurs.

654 The internal hydraulic jump is the only important feature observed in
 655 our simulations that was not mentioned by Baines (2008). It is possible,
 656 however, that the measurements performed in the laboratory experiments

were not adequate to detect this phenomenon. Such jumps are common features of oceanic outflows (e.g., Pratt et al., 2007; Legg et al., 2009), but their presence does not guarantee that flow splitting will occur. Instead, our results show that all the following characteristics are necessary for the development of flow splitting: $B_0 \sim 10^{-2}$, $R > 2$, and the development of an internal hydraulic jump.

Recent laboratory studies (Cortés et al., 2014b) and numerical simulations (Cortés et al., 2015) of gravity currents encountering a density step in a two-layered stratified basin have stressed the importance of the internal properties in determining the behavior of the flow. It was shown that gravity currents are more likely to split as the density interface becomes more diffusive. That is, the flow is more likely to detrain into two parts as Fr increases (equivalent to Ri decreasing). Our numerical experiments show that, in a linearly stratified environment where the density current is initially supercritical ($Fr > 1$), flow splitting is more likely to occur as Ri (Fr) increases (decreases). This disparity is likely a direct consequence of the different types of instabilities that may occur in these systems.

4.2. Application to oceanic outflows

We now compare α versus Ri_0 for all the numerical experiments conducted here (Fig. 15, see also Table 2 for a full list of all parameters). The experiments are separated into three groups, based on B_0 , and the different flow regimes are plotted with different colors. Regardless of α and Ri_0 , when $B_0 = 0.002$ the outflow always generates an underflow (Fig. 15a). Similarly, when $B_0 = 0.13$, the system always results in an interflow (Fig. 15c). However, when $B_0 = 0.016$, the type of regime is dictated by α and Ri_0 (Fig. 15b).

Table 2: Parameters for all numerical simulations conducted here (see text for the definition). The data is grouped based on Pr , with a) $Pr = 1$; b) $Pr = 10$; and c) $Pr = 50$.

	Q_0 ($\text{m}^2 \text{ s}^{-1}$)	g'_0 (m s^{-2})	N (s^{-1})	α	B_0	R	Ri_0	Re_e
a)	50.6	1.8×10^{-3}	5.0×10^{-4}	0.025	0.002	1.4	0.19	10.12
	50.4	1.8×10^{-3}	5.0×10^{-4}	0.05	0.002	1.3	0.17	10.08
	51.4	1.8×10^{-3}	5.0×10^{-4}	0.1	0.002	1.2	0.13	10.28
	47.8	1.7×10^{-3}	1.0×10^{-3}	0.025	0.016	1.6	0.21	9.56
	48.3	1.7×10^{-3}	1.0×10^{-3}	0.05	0.016	1.3	0.15	9.66
	48.2	1.7×10^{-3}	1.0×10^{-3}	0.1	0.016	1.2	0.13	9.64
	62.6	2.7×10^{-3}	2.5×10^{-3}	0.025	0.130	1.3	0.15	8.35
	62.1	2.7×10^{-3}	2.5×10^{-3}	0.05	0.130	1.2	0.14	8.28
	61.2	2.7×10^{-3}	2.5×10^{-3}	0.1	0.130	0.9	0.10	8.16
b)	51.2	1.8×10^{-3}	5.0×10^{-4}	0.025	0.002	2.3	0.27	10.24
	51.8	1.8×10^{-3}	5.0×10^{-4}	0.05	0.002	1.9	0.20	10.36
	50.7	1.8×10^{-3}	5.0×10^{-4}	0.1	0.002	1.8	0.14	10.14
	48.2	1.7×10^{-3}	1.0×10^{-3}	0.025	0.016	2.4	0.29	9.64
	48.0	1.7×10^{-3}	1.0×10^{-3}	0.05	0.016	2.1	0.21	9.60
	47.7	1.7×10^{-3}	1.0×10^{-3}	0.1	0.016	2.0	0.18	9.54
	62.5	2.7×10^{-3}	2.5×10^{-3}	0.025	0.130	2.3	0.30	8.33
	61.5	2.7×10^{-3}	2.5×10^{-3}	0.05	0.130	1.8	0.22	8.20
	61.2	2.7×10^{-3}	2.5×10^{-3}	0.1	0.137	1.8	0.15	8.16
c)	51.2	1.8×10^{-3}	5.0×10^{-4}	0.025	0.002	2.3	0.28	10.24
	51.8	1.8×10^{-3}	5.0×10^{-4}	0.05	0.002	2.2	0.22	10.36
	51.4	1.8×10^{-3}	5.0×10^{-4}	0.1	0.002	2.0	0.16	10.28
	48.5	1.7×10^{-3}	1.0×10^{-3}	0.025	0.016	2.7	0.31	9.70
	48.0	1.7×10^{-3}	1.0×10^{-3}	0.05	0.016	2.4	0.24	9.60
	48.1	1.7×10^{-3}	1.0×10^{-3}	0.1	0.016	2.3	0.20	9.62
	62.0	2.7×10^{-3}	2.5×10^{-3}	0.025	0.130	2.5	0.32	8.27
	62.6	2.7×10^{-3}	2.5×10^{-3}	0.05	0.130	2.2	0.24	8.35
	61.7	2.7×10^{-3}	2.5×10^{-3}	0.1	0.139	2.2	0.18	8.23

For $\alpha = 0.1$, split flows occur when $Ri_0 \geq 0.18$ and an interflow occurs otherwise. As the slope becomes shallower, the conditional Ri_0 required for flow splitting increases to 0.20 and 0.29, for $\alpha = 0.05$ and 0.025, respectively.

For comparison, Fig. 15 also shows the parameter values at the top of the slope (i.e., at the level of the initial outflow) inferred from observations of selected oceanic outflows (see Table 3). The B_0 values derived from observations of oceanic outflows are between 9.0×10^{-3} and 7.2×10^{-1} , while α varies between 0.03 and 0.15. The bulk Richardson number, J , was estimated by assuming $\Delta U \sim U$. Microstructure measurements of oceanic outflows are very rare and we could not infer h_u and R from the available observations. Therefore, estimates from observations are plotted by assuming $h_u \sim h_0/3$ and three values of R (1, 1.5 and 2.0). Lastly, following Baines (2005), the Reynolds number is estimated as $Re = \frac{Q_0}{\nu}$, where ν is the molecular viscosity ($\nu = 1.0 \times 10^{-6} \text{ m}^2 \text{ s}^{-1}$).

The parameters estimated for the Ross Sea outflows (orange stars in Fig. 15b) fall in the range of B_0 that, based on our numerical experiments, are propitious to either the interflow regime (when $Ri_0 = 0.1$) or the flow splitting regime (when $Ri_0 = 0.2$); $Ri_0 = 0.15$ falls between the two regimes, so it is not clear which regime occurs in this case. The parameters for the Denmark Straits, Mediterranean, Weddell, and Red Seas fall within numerical experiments that result in the interflow regime (Fig. 15c). None of the estimates based on observations of oceanic outflows fall within numerical experiments that generate the underflow regime. This is unexpected since observations (see references in Table 3) show that the outflows from the Denmark Straits, Weddell and Ross Seas sink to the bottom of their respective

Figure 15: Shelf slope (α) versus center Richardson number ($Ri_0 = J \times R$) for all the simulations conducted here (see also Table 2). The results are divided into three groups based on the initial buoyancy number (B_0): a) $B_0 < 8.0 \times 10^{-3}$; b) $8.0 \times 10^{-3} \leq B_0 \leq 8.0 \times 10^{-2}$; c) $B_0 > 8.0 \times 10^{-2}$. Simulated underflows, split flows and interflows, are represented by the blue, red and green squares, respectively. Also shown are the values for various oceanic outflows, representing the conditions at the top of the slope, inferred from available observations (stars, see Table 3 for additional information).

ocean basins. It is possible that two factors that are likely to occur in oceanic outflows but are not considered in our analyses may contribute to the descent of these outflows. First, the parameters derived from observations shown in Table 3 and Fig. 15 represent the conditions at the top of the shelf slope and these are likely to vary significantly further downslope, potentially leading to transitions between different regimes. In particular, the ambient stratification (N) in these outflows is likely to be significantly smaller at larger depths. Second, the thermobaric effect (Gill, 1973; Matsumura and Hasumi, 2010) might be important in maintaining the density anomaly as the outflow reaches large depths.

Figure 15 does not include the results from the two laboratory experiments presented by Baines (2005, 2008) that generated split flows, since the parameter values in these experiments differ significantly from the range investigated in our numerical simulations and observations; however, these values are listed in Table 3. To generate a turbulent flow (i.e., large enough Re values), the slope used in these experiments was much larger than those observed in nature and this is a general drawback of laboratory experiments of turbulent gravity currents. In addition, the ambient stratification frequency used in these experiments were at least two orders of magnitude larger than the typical values next to the major oceanic outflows (Table 3). Both parameters (α and N) are important in setting the conditions under which different turbulent regimes may occur. Therefore, we conclude that the flow regime criteria obtained in laboratory experiments of dense fluid down slope and into stratified environments differ from that obtained under geophysical parameters and scales. Based on our results and observations taken in the

Table 3: Parameters representing the conditions at the top of the slope for selected oceanic outflows estimated based on the observations presented in the following articles: Ross Sea: Gordon et al. (2004); Muench et al. (2009); Mediterranean Sea: Baringer and Price (1997); Price et al. (1993); Weddell Sea: Foldvik et al. (2004); Denmark Straits: Girtton and Sanford (2003); Red Sea: Matt and Johns (2007); Peters et al. (2005). The parameters for the two laboratory experiments conducted by Baines (2005) that resulted in flow splitting are also shown.

Region	α	g'_0 (m s ⁻²)	N (s ⁻¹)	h_0 (m)	U_0 (m s ⁻¹)	Q_0 (m ² s ⁻¹)	B_0	J	Re
Ross Sea	0.11	2×10^{-3}	7.0×10^{-4}	150	1.0	150	1.3×10^{-2}	0.10	1.5×10^8
Med. Sea	0.15	1.6×10^{-3}	2.0×10^{-3}	200	1.0	200	7.2×10^{-1}	0.11	2.0×10^8
Red Sea	0.004	1.4×10^{-2}	7.0×10^{-3}	100	0.55	55	9.6×10^{-2}	0.19	5.5×10^7
Weddell Sea	0.03	2×10^{-3}	7.0×10^{-4}	100	1.0	100	9.0×10^{-3}	0.07	1.0×10^8
Denmark Straits	0.03	2×10^{-3}	1.3×10^{-3}	200	0.7	140	7.7×10^{-2}	0.27	1.4×10^8
Baines	0.72	25.01	1.22	0.27	2.27	0.61	1.8×10^{-3}	1.13	61.4
Baines	0.72	22.65	1.16	0.61	2.16	1.32	4.0×10^{-3}	2.55	132.0

732 Ross and Weddell Seas, we hypothesize that flow splitting around Antarc-
733 tica could carry water properties from the shelf into the deep ocean at two
734 distinct depths.

735 4.3. Influence of the maximum ocean depth

736 So far, the influence of the maximum ocean depth (H_{max}) on the develop-
737 ment of the outflow has been ignored. Assuming conditions where an outflow
738 would equilibrate at a depth H_e (i.e., become neutrally buoyant), a very deep
739 ocean ($H_{max} \gg H_e$) would never result in an underflow. Similarly, a very
740 shallow ocean ($H_{max} \ll H_e$) would never produce an interflow. It is there-
741 fore useful to quantify how the vertical distance between the top (H_{shelf}) and

742 bottom (H_{max}) of the slope ($\Delta H = H_{max} - H_{shelf}$) compares with the neu-
 743 tral buoyancy level ($H_e = g'_0/N^2$), relative to H_{shelf} , that the outflow would
 744 achieve in the absence of mixing. The relevant non-dimensional parameter,
 745 $H^* = \Delta H/H_e$, is expected to be small for underflows and large for interflows.
 746 In our numerical experiments, where ΔH is always 1500 m, H^* is 0.21, 0.88
 747 and 3.47 for cases where underflows, split flows and interflows are generated,
 748 respectively (i.e., when $B_0 = 0.002$, 0.016 and 0.13, respectively; see Table 3).
 749 Since we kept ΔH constant, changing B_0 by varying N and/or g'_0 is equiva-
 750 lent to changing H^* . However, it is possible to change H^* while keeping B_0
 751 fixed (i.e., by changing ΔH) and this could modify the flow regimes obtained
 752 based only on B_0 . For example, it is possible that if we were to repeat the set
 753 of experiments where $B_0 = 0.002$ using a larger ΔH the flow would result in
 754 an interflow or a split flow rather than an underflow. Therefore, although we
 755 have not varied H_0 in our set of experiments, the influence of this parameter
 756 should also be considered when inferring the flow regime of the system.

757 For the real oceanic outflows listed in Table 3, $H^* \gg 1$ for the Mediter-
 758 ranean Sea, the Red Sea and the Denmark Straits outflows, suggesting in-
 759 terflows (where $\Delta H \sim 3000$, 1500, 3500 m, respectively), while $H^* \sim 0.63$
 760 for the Ross Sea outflow ($\Delta H \sim 2500$ m) and 0.88 for Weddell Sea outflow
 761 ($\Delta H \sim 3500$ m), suggesting flow splitting. This classification is consistent
 762 with the regimes proposed based on B_0 .

763 4.4. *Assumptions, limitations and recommendations for future studies*

764 The results discussed thus far are all based on 2D simulations. In terms
 765 of turbulence, there is a fundamental difference between 2D and 3D, where
 766 energy cascades towards larger scales in the first and towards smaller scales

767 in the latter (e.g., Vallis, 2006). Since one of the goals of this study is to
 768 investigate the conditions under which the flow splitting regime may occur in
 769 oceanic outflows, this distinction puts in doubt whether this regime can be
 770 generated in a 3D system. In the context of large-scale oceanic outflows, the
 771 numerical simulations conducted by Özgökmen et al. (2004) show that mixing
 772 in 2D simulations is overestimated when compared to the values obtained
 773 in laboratory experiments and in 3D simulations. The authors attributed
 774 this difference to the absence of secondary instabilities in 2D, while in 3D
 775 such instabilities develop in the along-slope direction and lead to a rapid
 776 breakdown of the vortices. Numerical simulations of gravity-current fronts in
 777 the lock-exchange configuration also revealed that vortices in 2D are unable
 778 to break down as they do in 3D, which therefore affects mixing and the
 779 density field (Härtel et al., 2000). Although it is likely that flow splitting
 780 will also occur in 3D simulations, we anticipate that in fully 3D systems the
 781 transition from one regime to another and the correspondent value of B_0 for
 782 each regime may differ from those presented here.

783 In addition to three-dimensionality, most oceanic outflows are also strongly
 784 influenced by rotation. For unconstrained flows going down broad slopes,
 785 which is the case of our simulations, rotation will balance the downward flow
 786 and therefore result in a large portion of the dense fluid following f/H con-
 787 tours, where f is the Coriolis parameter and H is the total water depth.
 788 In this case, an additional flow regime where the plume is arrested can also
 789 occur (Wobus et al., 2013). In terms of mixing, the rotating experiments of
 790 unconstrained flows reported by Wells (2007) show entrainment rates similar
 791 to those obtained in similar non-rotating experiments reported by Ellison and

792 Turner (1959). However, in the presence of vertical walls or inclined bound-
 793 aries, e.g., ridges or canyons, secondary flows induced by rotation lead to
 794 more vigorous entrainment than in non-rotating systems. Topographic fea-
 795 tures are often present along continental slopes around the oceans, including
 796 regions of major oceanic outflows (e.g., see Fig.1 in Allen and Durrieu de
 797 Madron, 2009). Since submarine canyons and ridges increase entrainment
 798 rates (Wåhlin et al., 2008), the transition B_0 values for the different flow
 799 regimes are likely to differ from those presented here. Therefore, future re-
 800 search should consider the effects of both rotation and three-dimensionality,
 801 ideally in the presence of small-scale topographic features, to confirm if the
 802 outcomes of this study are still valid under such conditions.

803 The low effective Reynolds number ($Re_e = Q_0/\nu_h$) values of the outflow
 804 in our numerical experiments (see Table 2) are also a limitation that needs
 805 to be addressed. Although these values are similar (i.e., same order of mag-
 806 nitude) to the Re values obtained in the laboratory experiment conducted
 807 by Baines (2005, 2008), they are much smaller than Re values observed next
 808 to the major oceanic outflows (Table 3). Thus, future numerical studies are
 809 required to confirm if the outcomes of this study also hold under larger effec-
 810 tive Reynolds number. This will require the application of direct numerical
 811 simulations (DNS), where all scales of motion down to the Kolmogorov scale
 812 are explicitly resolved (i.e., without any turbulence closure). Although this
 813 method is computationally very expensive (even for a 2D setup similar to
 814 the one employed here), the results presented here can be used as a guide to
 815 narrow the parameter space to be explored.

816 Finally, in all simulations conducted here, the outflow is initially super-

critical ($Fr > 1$) since most oceanic outflows fall in this category (Legg et al., 2009). To simulate subcritical conditions following our choice of parameters and scales would require much higher vertical resolution, since subcritical outflows, in general, tend to develop step-like velocity and density profiles (e.g., Cortés et al., 2014b). Hence, whether the outcomes of this study apply to subcritical conditions remains to be tested. Nevertheless, we do not expect that the flow splitting regime will occur in subcritical outflows since internal hydraulic jumps are not likely to develop in such conditions.

5. Concluding remarks

A comparison between our numerical results and oceanic observations suggests that flow splitting may occur in dense-water outflows with weak ambient stratification, such as Antarctic outflows. Oceanic outflows, in general, are not well represented in the current generation of climate models and this is one of the main challenges for improving climate predictions (e.g., Heuzé et al., 2013; Legg et al., 2009). Therefore, it would be extremely valuable to confirm whether flow splitting occurs in these systems. The outcomes of the present study can be used in the design of observational programs that aim to study oceanic outflows, as well as in future numerical and laboratory studies. Perhaps the best way to observe flow splitting in these systems would involve the release of some kind of tracer (e.g., rhodamine) into the dense water before it flows downslope. Although this procedure may not be feasible in large-scale outflows, given the extremely large amount of dye that would be required and the long time needed to observe its offshore response, a dye release into a smaller scale system such as a fjord overflow would be

841 much more feasible. Anthropogenic tracers, such as radionuclide and chlo-
842 rofluorocarbons, may provide a possible alternative to detect flow splitting
843 in large-scale outflows.

844 Finestructure and microstructure measurements of velocity and buoy-
845 ancy profiles within outflows would provide valuable support for the pro-
846 cesses suggested by our simulations. This information can then be used, in
847 combination with other environmental parameters (i.e., ambient stratifica-
848 tion and bottom slope), to identify the flow regimes based on the results
849 presented here. However, such comparison must be performed with caution
850 since three-dimensionality and rotation (both absent in our simulations) will
851 affect mixing and, therefore, the values representing the transition from each
852 flow regime are likely to differ from those obtained based on our numerical
853 experiments. In addition, the parameters governing the flow regimes are ex-
854 pected to vary downstream of the initial level of the outflow, perhaps leading
855 to a transition between different flow regimes.

856 Lastly, for simplicity and lack of knowledge on what value Pr takes in real
857 oceanic outflows, numerical studies often assume $Pr = 1$ (Özgökmen et al.,
858 2004, 2006; Özgökmen and Fischer, 2008; Wobus et al., 2013). Our study
859 demonstrates that flow splitting occurs when $Pr \geq 10$ and this highlights
860 the need to further investigate the actual form of Pr in oceanic outflows so
861 that appropriate values can be used in numerical models.

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869 References

- 870 Aagaard, K., Swift, J., Carmack, E., 1985. Thermohaline circulation in the
871 arctic mediterranean seas. *J. Geophys. Res.* 90, 4833–4846.
- 872 Adcroft, A., Hill, C., Marshall, J., 1997. Representation of topography by
873 shaved cells in a height coordinate ocean model. *Mon. Weather Rev.* 125,
874 2293–2315.
- 875 Allen, S., Durrieu de Madron, X., 2009. A review of the role of submarine
876 canyons in deep-ocean exchange with the shelf. *Ocean Science* 5, 607–620.
- 877 Baines, P.G., 2001. Mixing in flows down gentle slopes into stratified envi-
878 ronments. *J. Fluid Mech.* 443, 237–270.
- 879 Baines, P.G., 2005. Mixing regimes for the flow of dense fluid down slopes
880 into stratified environments. *J. Fluid Mech.* 538, 245–267.
- 881 Baines, P.G., 2008. Mixing in downslope flows in the ocean-plumes versus
882 gravity currents. *Atmosphere-Ocean* 46, 405–419.
- 883 Baringer, M., Price, J., 1997. Mixing and spreading of the Mediterranean
884 outflow. *J. Phys. Oceanogr.* 27, 1654–1677.
- 885 Cenedese, C., Adduce, C., 2008. Mixing in a density-driven current flowing
886 down a slope in a rotating fluid. *J. Fluid Mech.* 604, 369–388.
- 887 Cenedese, C., Whitehead, J.A., Ascarelli, T., Ohiwa, M., 2004. A dense cur-
888 rent flowing down a sloping bottom in a rotating fluid. *J. Phys. Oceanogr.*
889 34, 188–203.

890 Cortés, A., Fleenor, W.E., Wells, M.G., de Vicente, I., Rueda, F.J., 2014a.
891 Pathways of river water to the surface layers of stratified reservoirs. *Lim-*
892 *nology and Oceanography* 59, 233–250.

893 Cortés, A., Rueda, F.J., Wells, M.G., 2014b. Experimental observations of
894 the splitting of a gravity current at a density step in a stratified water
895 body. *J. Geophys. Res.* 119, 1038–1053.

896 Cortés, A., Wells, M.G., Fringer, O.B., Arthur, R., Rueda, F.J., 2015. Nu-
897 merical investigation of split flows by gravity currents into two-layered
898 stratified water bodies. *J. Geophys. Res.* 120, 5254–5271.

899 De Cesare, G., Boillat, J.L., Schleiss, A.J., 2006. Circulation in stratified
900 lakes due to flood-induced turbidity currents. *J. Environ. Eng.* 132, 1508–
901 1517.

902 Elliott, Z.A., Venayagamoorthy, S.K., 2011. Evaluation of turbulent Prandtl
903 (Schmidt) number parameterizations for stably stratified environmental
904 flows. *Dyn. Atmos. Oceans* 51, 137–150.

905 Ellison, T.H., Turner, J.S., 1959. Turbulent entrainment in stratified flows.
906 *J. Fluid Mech.* 6, 423–448.

907 Ezer, T., 2005. Entrainment, diapycnal mixing and transport in three-
908 dimensional bottom gravity current simulations using the mellor–yamada
909 turbulence scheme. *Ocean Modell.* 9, 151–168.

910 Foldvik, A., Gammelsrød, T., Østerhus, S., Fahrbach, E., Rohardt, G.,
911 Schröder, M., Nicholls, K.W., Padman, L., Woodgate, R.A., 2004. Ice

912 shelf water overflow and bottom water formation in the southern Weddell
913 Sea. *J. Geophys. Res.* 109, 1–15.

914 Gill, A., 1973. Circulation and bottom water production in the Weddell Sea.
915 *Deep Sea Res.* 20, 111–140.

916 Girton, J.B., Sanford, T.B., 2003. Descent and modification of the overflow
917 plume in the Denmark strait*. *J. Phys. Oceanogr.* 33, 1351–1364.

918 Gordon, A., Zambianchi, E., Orsi, A., 2004. Energetic plumes over the
919 western Ross Sea continental slope. *Geophys. Res. Letters* 31.

920 Griffies, S.M., Pacanowski, R.C., Hallberg, R.W., 2000. Spurious diapycnal
921 mixing associated with advection in a z-coordinate ocean model. *Mon.*
922 *Weather Rev.* 128, 538–564.

923 Guo, Y., Zhang, Z., Shi, B., 2014. Numerical Simulation of Gravity Current
924 Descending a Slope into a Linearly Stratified Environment. *J. Hydraulic*
925 *Eng.* 140.

926 Härtel, C., Meiburg, E., Necker, F., 2000. Analysis and direct numerical
927 simulation of the flow at a gravity-current head. part 1. flow topology and
928 front speed for slip and no-slip boundaries. *J. Fluid Mech.* 418, 189–212.

929 Heuzé, C., Heywood, K.J., Stevens, D.P., Ridley, J.K., 2013. Southern Ocean
930 bottom water characteristics in CMIP5 models. *Geophys. Res. Letters* 40,
931 1409–1414.

932 Hogg, A.M., Ivey, G., 2003. The Kelvin-Helmholtz to Holmboe instability
933 transition in stratified exchange flows. *J. Fluid Mech.* 477, 339–362.

934 Holland, D.M., Rosales, R.R., Stefanica, D., Tabak, E.G., 2002. Internal
935 hydraulic jumps and mixing in two-layer flows. *J. Fluid Mech.* 470, 63–83.

936 Howard, L.N., 1961. Note on a paper of John W. Miles. *J. Fluid Mech.* 10,
937 509–512.

938 Hughes, G., Griffiths, R., 2006. A simple convective model of the global over-
939 turning circulation, including effects of entrainment into sinking regions.
940 *Ocean Modell.* 12, 46–79.

941 Ilicak, M., Adcroft, A.J., Griffies, S.M., Hallberg, R.W., 2012. Spurious
942 dianeutral mixing and the role of momentum closure. *Ocean Modell.* 45-
943 46, 37–58.

944 Jackson, L., Hallberg, R., Legg, S., 2008. A parameterization of shear-driven
945 turbulence for ocean climate models. *J. Phys. Oceanogr.* 38, 1033–1053.

946 Lane-Serff, G., Baines, P., 1998. Eddy formation by dense flows on slopes in
947 a rotating fluid. *J. Fluid Mech.* 363, 229–252.

948 Legg, S., Ezer, T., Jackson, L., Briegleb, B., Danabasoglu, G., Large, W.,
949 Wu, W., Chang, Y., Özgökmen, T.M., Peters, H., Xu, X., Chassignet,
950 E.P., Gordon, A.L., Griffies, S., Hallberg, R., Price, J., Riemenschneider,
951 U., Yang, J., 2009. Improving oceanic overflow representation in climate
952 models: the gravity current entrainment climate process team. *Bulletin of*
953 *the American Meteorological Society* 90, 657–670.

954 Legg, S., Hallberg, R., Girton, J., 2006. Comparison of entrainment in
955 overflows simulated by σ -coordinate, isopycnal and non-hydrostatic mod-
956 els. *Ocean Modell.* 11, 69–97.

- 957 Legg, S., Klymak, J., 2008. Internal hydraulic jumps and overturning gener-
958 ated by tidal flow over a tall steep ridge. *J. Phys. Oceanogr.* 38, 1949–1964.
- 959 Long, R.R., 1953. Some aspects of the flow of stratified fluids: I. A theoretical
960 investigation. *Tellus* 5, 42–58.
- 961 Long, R.R., 1954. Some aspects of the flow of stratified fluids: II. Experi-
962 ments with a two-fluid system. *Tellus* 6, 97–115.
- 963 Marques, G.M., Özgökmen, T.M., 2014. On modeling turbulent exchange in
964 buoyancy-driven fronts. *Ocean Modell.* 83, 43–62.
- 965 Marshall, J., Adcroft, A., Hill, C., Perelman, L., Heisey, C., 1997a. A finite-
966 volume, incompressible navier stokes model for studies of the ocean on
967 parallel computers. *J. Geophys. Res.* 102, 5753–5766.
- 968 Marshall, J., Hill, C., Perelman, L., Adcroft, A., 1997b. Hydrostatic, quasi-
969 hydrostatic, and nonhydrostatic ocean modeling. *J. Geophys. Res.* 102,
970 5733–5752.
- 971 Matsumura, Y., Hasumi, H., 2010. Modeling ice shelf water overflow and
972 bottom water formation in the southern Weddell Sea. *J. Geophys. Res.*
973 115, 1–17.
- 974 Matt, S., Johns, W.E., 2007. Transport and entrainment in the Red Sea
975 outflow plume. *J. Phys. Oceanogr.* 37, 819–836.
- 976 Mauritzen, C., Price, J., Sanford, T., Torres, D., 2005. Circulation and
977 mixing in the Faroese channels. *Deep Sea Res. I* 52, 883–913.

- 978 McDougall, T.J., 1987. Thermobaricity, cabbeling, and water-mass conver-
979 sion. *J. Geophys. Res.* 92, 5448–5464.
- 980 Miles, J.W., 1961. On the stability of heterogeneous shear flows. *J. Fluid*
981 *Mech.* 10, 496–508.
- 982 Mitsudera, H., Baines, P.G., 1992. Downslope gravity currents in a continu-
983 ously stratified environment: a model of the bass strait outflow, in: 11th
984 Australasian Fluid Mechanics Conf., pp. 1017–1020.
- 985 Monaghan, J.J., 2007. Gravity current interaction with interfaces. *Annu.*
986 *Rev. Fluid Mech.* 39, 245–261.
- 987 Muench, R., Padman, L., Gordon, A., Orsi, A., 2009. A dense water outflow
988 from the Ross Sea, Antarctica: Mixing and the contribution of tides. *J.*
989 *Mar. Sys.* 77, 369–387.
- 990 Necker, F., Härtel, C., Kleiser, L., Meiburg, E., 2002. High-resolution sim-
991 ulations of particle-driven gravity currents. *Int. J. Multiphase Flow* 28,
992 279–300.
- 993 Özgökmen, T.M., Chassignet, E.P., 2002. Dynamics of two-dimensional tur-
994 bulent bottom gravity currents. *J. Phys. Oceanogr.* 32, 1460–1478.
- 995 Özgökmen, T.M., Fischer, P.F., 2008. On the role of bottom roughness in
996 overflows. *Ocean Modell.* 20, 336–361.
- 997 Özgökmen, T.M., Fischer, P.F., Duan, J., Iliescu, T., 2004. Three-
998 dimensional turbulent bottom density currents from a high-order nonhy-
999 drostatic spectral element model. *J. Phys. Oceanogr.* 34, 2006–2026.

- 1000 Özgökmen, T.M., Fischer, P.F., Johns, W.E., 2006. Product water mass
1001 formation by turbulent density currents from a high-order nonhydrostatic
1002 spectral element model. *Ocean Modell.* 12, 237–267.
- 1003 Padman, L., Howard, S.L., Orsi, A.H., Muench, R.D., 2009. Tides of the
1004 northwestern Ross Sea and their impact on dense outflows of Antarctic
1005 Bottom Water. *Deep Sea Res. I* 56, 818–834.
- 1006 Peters, H., Johns, W.E., Bower, A.S., Fratantoni, D.M., 2005. Mixing and
1007 entrainment in the Red Sea outflow plume. Part I: Plume structure. *J.*
1008 *Phys. Oceanogr.* 35, 569–583.
- 1009 Pratt, L.J., Riemenschneider, U., Helfrich, K.R., 2007. A transverse hy-
1010 draulic jump in a model of the Faroe Bank Channel outflow. *Ocean Modell.*
1011 19, 1–9.
- 1012 Price, J., Baringer, M., 1994. Outflows and deep water production by
1013 marginal seas. *Progr. in Oceanogr.* 33, 161–200.
- 1014 Price, J.F., Baringer, M.O., Lueck, R.G., Johnson, G.C., Ambar, I., Par-
1015 rilla, G., Cantos, A., Kennelly, M.A., Sanford, T.B., 1993. Mediterranean
1016 outflow mixing and dynamics. *Science* 259, 1277–1282.
- 1017 Simpson, J.E., 1999. Gravity currents: In the environment and the labora-
1018 tory. Cambridge University Press.
- 1019 Smyth, W., Carpenter, J., Lawrence, G., 2007. Mixing in symmetric holmboe
1020 waves. *J. Phys. Oceanogr.* 37, 1566–1583.

1021 Smyth, W., Winters, K., 2003. Turbulence and mixing in holmboe waves. *J.*
1022 *Phys. Oceanogr.* 33, 694–711.

1023 Strang, E., Fernando, H., 2001. Vertical mixing and transports through a
1024 stratified shear layer. *J. Phys. Oceanogr.* 31, 2026–2048.

1025 Vallis, G.K., 2006. Atmospheric and oceanic fluid dynamics: fundamentals
1026 and large-scale circulation. Cambridge University Press.

1027 Wåhlin, A., Cenedese, C., 2006. How entraining density currents influence
1028 the stratification in a one-dimensional ocean basin. *Deep Sea Res. II* 53,
1029 172–193.

1030 Wåhlin, A., Darelius, E., Cenedese, C., Lane-Serff, G., 2008. Laboratory
1031 observations of enhanced entrainment in dense overflows in the presence
1032 of submarine canyons and ridges. *Deep Sea Res. I* 55, 737–750.

1033 Wells, M., 2007. Influence of coriolis forces on turbidity currents and sediment
1034 deposition, in: *Particle-Laden Flow*. Springer, pp. 331–343.

1035 Wells, M., Nadarajah, P., 2009. The intrusion depth of density currents
1036 flowing into stratified water bodies. *J. Phys. Oceanogr.* 39, 1935–1947.

1037 Wells, M.G., Wettlaufer, J.S., 2007. The long-term circulation driven by
1038 density currents in a two-layer stratified basin. *J. Fluid Mech.* 572, 37–58.

1039 Winton, M., Hallberg, R., Gnanadesikan, A., 1998. Simulation of density-
1040 driven frictional downslope flow in z-coordinate ocean models. *J. Phys.*
1041 *Oceanogr.* 28, 2163–2174.

1042 Wobus, F., Shapiro, G.I., Huthnance, J.M., Maqueda, M.A.M., 2013. The
1043 piercing of the Atlantic Layer by an Arctic shelf water cascade in an ide-
1044 alised study inspired by the Storfjorden overflow in Svalbard. *Ocean Mod-*
1045 *ell.* 71, 54–65.