



RESEARCH ARTICLE

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Forcing of the overturning circulation across a circumpolar channel by internal wave breaking

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Key Points:

- The breaking of topographically generated internal lee waves enhances diapycnal diffusivity
- The diapycnal mixing generated by the simulated eddy field regulates the deep overturning
- Mesoscale eddies displace isopycnals vertically and set the vertical structure of the deep MOC

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Abstract The hypothesis that the impingement of mesoscale eddy flows on small-scale topography regulates diapycnal mixing and meridional overturning across the deep Southern Ocean is assessed in an idealized model. The model simulates an eddying circumpolar current coupled to a double-celled meridional overturning with properties broadly resembling those of the Southern Ocean circulation and represents lee wave-induced diapycnal mixing using an online formulation grounded on wave radiation theory. The diapycnal mixing generated by the simulated eddy field is found to play a major role in sustaining the lower overturning cell in the model, and to underpin a significant sensitivity of this cell to wind forcing. The vertical structure of lower overturning is set by mesoscale eddies, which propagate the effects of near-bottom diapycnal mixing by displacing isopycnals vertically.

1. Introduction

The Southern Ocean plays a pivotal role in the global overturning circulation. The absence of continental barriers in the latitude band of Drake Passage permits the existence of the Antarctic Circumpolar Current (ACC), which is supported geostrophically by sloping isopycnals and serves as the main conduit for oceanic exchanges between the three major ocean basins. Coupled to this eastward flow, a meridional circulation exists in which deep water of predominantly North Atlantic origin together with Indian and Pacific Deep waters upwells along the poleward-shoaling isopycnals of the ACC and is returned equatorward in a double overturning cell [see *Rintoul and Garabato*, 2013; *Talley*, 2013, and references therein]. Our present understanding of the dynamics regulating these zonal and meridional circulations and their coupling owes much to studies of idealized models of (or including) a zonally reentrant channel, and to residual-mean analyses of realistic models [e.g., *Marshall and Radko*, 2003; *Hallberg and Gnanadesikan*, 2006; *Abernathey et al.*, 2011; *Dufour et al.*, 2012; *Meredith et al.*, 2012; *Marshall and Speer*, 2012; *Morrison and Hogg*, 2013; *Thompson and Naveira Garabato*, 2014]. This literature unanimously emphasizes the significance of mesoscale eddies, generated by baroclinic instability of the ACC, in both closing the zonal momentum budget and inducing meridional overturning, particularly at open-channel depths where no time-mean geostrophic flow may exist.

More recently, evidence has emerged suggesting that the Southern Ocean eddy field may impact meridional overturning beyond its widely acknowledged role in effecting adiabatic, along-isopycnal transport. This evidence comes from two distinct lines of work. First, theoretical estimates of the rate of generation of internal lee waves by eddy flows impinging on small-horizontal-scale ($O(1-10\text{ km})$) topography indicate that this mechanism may be a significant contributor to the dissipation of the eddy field [*Nikurashin and Ferrari*, 2010a, 2010b; *Scott et al.*, 2011]. Second, fine and microstructure measurements across Southern Ocean areas of complex topography reveal elevated (typically by at least 1 order of magnitude relative to oceanic background values) rates of turbulent dissipation and diapycnal mixing in the deepest 1000–2000 m of those regions, linked to the eddy-induced generation and breaking of lee waves [*Naveira Garabato et al.*, 2004; *St. Laurent et al.*, 2012; *Breareley et al.*, 2013; *Sheen et al.*, 2013, 2014; *Waterman et al.*, 2013]. As diapycnal mixing has been shown to be a rate-setting process in deep-ocean overturning [e.g., *Ito and Marshall*, 2008; *Nikurashin and Vallis*, 2011], and eddy kinetic energy in the ACC to be highly sensitive to wind perturbations [see *Meredith et al.*, 2012, and references therein], the hypothesis is prompted that the eddy field's impingement on small-scale topography (and, by extension, wind forcing) regulates diapycnal mixing and meridional overturning across the deep Southern Ocean.

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In this paper, we assess the validity of this hypothesis in an idealized model of an eddying zonal channel flow coupled to a double-celled meridional overturning with properties broadly resembling those of the Southern Ocean circulation. Lee wave-induced diapycnal mixing is represented in the model using an online formulation grounded on wave radiation theory. We show that the lee wave-induced diapycnal mixing generated by the simulated eddy field is a major factor in sustaining the lower overturning cell in the model, and that it underpins a significant sensitivity of this cell to wind forcing perturbations.

2. Parameterization of Lee Wave-Induced Diapycnal Mixing

Our parameterization of the spatiotemporally variable diapycnal diffusivity due to lee wave breaking, κ_{iw} , takes the form

$$\kappa_{iw}(x, y, z, t)N^2(x, y, z, t) = -B_{iw}(x, y, z, t) = \rho_0^{-1}\Gamma G(x, y, t)F(z), \quad (1)$$

where N is the buoyancy frequency, $B_{iw}(x, y, z, t)$ is the turbulent buoyancy flux driven by breaking internal waves, ρ_0 is density, $\Gamma = 0.2$ is the mixing efficiency for shear-driven turbulence [Osborn, 1980]; G is the rate of energy transfer from mesoscale eddy motions, resolved in the model, to the unresolved lee wave field; and F is a vertical structure function of the rate of IW energy dissipation. G is expressed as a function of near-bottom horizontal velocity and N , and involves a 2-D spectral description of the unresolved small-scale bottom topography implicated in lee wave generation. Its specification, included below for completeness, follows the wave radiation theory-based formulation of Scott *et al.* [2011].

We have chosen the canonical value for mixing efficiency for our setup as postulated by Osborn [1980]. In a recent study, de Lavergne *et al.* [2015] have explored the impact of variable mixing efficiency on the deep overturning circulation and found that a mixing efficiency corresponding to $\Gamma = 0.2$ is a reasonable assumption, although it may overestimate the near-bottom eddy diffusivities due to the weak stratification near the bottom boundary.

In this work, we will assume that the lee waves generated by geostrophic flow over topography do not experience significant horizontal propagation (relative to the model grid spacing) and that they break at some level within the water column, such that the depth-integrated turbulent kinetic energy dissipation (ϵ) matches the vertical flux of energy associated with the waves at generation (G), i.e.

$$\int_{-H}^0 \epsilon(x, y, z, t) dz = G(x, y, t). \quad (2)$$

The first assumption follows from the intrinsic role that the geostrophic flow plays in the generation and persistence of the lee waves, which locks them to the horizontal proximity of their generation site [e.g., Ferrari and Nikurashin, 2010]. The second reflects our present lack of knowledge on the characteristic evolution of lee waves as they propagate vertically, and the high likelihood of eventual breaking, particularly in the presence of nonlinearity in the internal wave field and buoyancy scaling effects (see Naveira Garabato *et al.* [2004] and Nikurashin and Ferrari [2010a] for a discussion). Note, though, that if the dissipation of the lee waves within the bottom part of the water column is not complete, the depth integral of ϵ will be reduced proportionally.

The transfer of energy into the internal lee waves, G , is of course associated [cf. Naveira Garabato *et al.*, 2013] with a loss of energy from the resolved flow via a wave drag $-\tau_{iw}$, consistent with $-\mathbf{u} \cdot \tau_{iw} = G$. We did not choose to include this in the following numerical simulations since we wished to consider the dynamic effect of the IW mixing in isolation.

2.1. Parameterization of Lee Wave Power

The generation of internal lee waves by geostrophic flow over topography has been studied from a theoretical perspective by Bell [1975], Gill [1982], and Nikurashin and Ferrari [2010a], among others. Adopting the traditional approximation of ignoring the horizontal components of the Coriolis force, the linearized density equation and nonhydrostatic, vertical momentum equation provide a relation between pressure and vertical velocity. For small-amplitude sinusoidal topography, the vertical velocity is imposed by the kinematic boundary condition, and one can solve for the vertical flux of energy resulting from the pressure-topographic slope correlation. The energy flux is nonzero if the oscillations generated by flow over

topography can propagate as waves, i.e., if their intrinsic frequency $\sigma = \mathbf{k} \cdot \mathbf{u}$ (where $\mathbf{k} = (k, l)$ is the horizontal wave number of the topography and $\mathbf{u} = (u, v)$ is the near-bottom horizontal velocity) lies between the inertial (f) and buoyancy (N) frequencies, $|f| < |\sigma| < N$. The energy extracted from the background flow is carried upward by the waves as they radiate away from the boundary, and dissipated when the waves break. Representing the topography by a two-dimensional power spectrum $P(k, l)$ that varies slowly on spatial scales much larger than the wavelengths of lee waves, it may be shown that the net vertical flux of energy resulting from the sum of all contributions within the internal wave band is given by

$$G_{\text{lin}}(x, y, t) = \frac{\rho_0}{4\pi^2} \int_{-\infty}^{\infty} \int_{-\infty}^{\infty} S(\sigma) \frac{\sigma}{\sqrt{k^2 + l^2}} P(k, l) \sqrt{N^2 - \sigma^2} \sqrt{\sigma^2 - f^2} dk dl, \tag{3}$$

where ρ_0 is the fluid density and

$$S(\sigma) = \begin{cases} 1 & \text{if } |f| < |\sigma| < N; \sigma > 0, \\ -1 & \text{if } |f| < |\sigma| < N; \sigma < 0, \\ 0 & \text{otherwise.} \end{cases}$$

The linear theory above describes well the generation of lee waves when the topography has small amplitude in the sense of

$$\frac{1}{\text{Fr}} \equiv \frac{\mathcal{H}N}{\mathcal{U}} \leq \frac{1}{\text{Fr}_c}, \tag{4}$$

where Fr is the Froude number of the flow, \mathcal{U} is a background velocity scale, \mathcal{H} is the amplitude of the topographic variations, and Fr_c is a critical Froude number. As topographic amplitude and buoyancy frequency increase or as background velocity decreases, the flow becomes increasingly blocked by the topography, leading to substantially less lee wave radiation than indicated by (3). Empirical corrections for this phenomenon have been put forward and tested with laboratory experiments and numerical simulations. These are reviewed by Scott *et al.* [2011] and will not be discussed here. We follow Scott *et al.* [2011] in adopting a correction for topographic blocking of the form

$$G(x, y, t) = G_{\text{lin}}(x, y, t) \frac{\mathcal{L}}{\pi} \left(\arccos(1 - 2\mathcal{L}) - 2(1 - 2\mathcal{L})\sqrt{\mathcal{L}(1 - \mathcal{L})} \right), \tag{5}$$

where

$$\begin{aligned} \mathcal{L} &= 1, & \text{Fr}^{-1} &\leq \text{Fr}_c^{-1}; \\ \mathcal{L} &= \frac{\text{Fr}}{\text{Fr}_c}, & \text{Fr}^{-1} &> \text{Fr}_c^{-1}. \end{aligned}$$

As in Scott *et al.* [2011], Fr_c^{-1} is taken as 0.7, and the amplitude of the topography in (5) is defined as

$$\mathcal{H}^2 = \frac{1}{4\pi^2} \int_{-\infty}^{\infty} \int_{-\infty}^{\infty} |S(\sigma)| P(k, l) dk dl. \tag{6}$$

This is a metric of the amplitude of the topographic variations contributing to lee wave generation and is therefore time dependent (through σ). Our results are insensitive (to better than 1%) to the choice of other possible corrections in the literature. See Naveira Garabato *et al.* [2013] for a discussion of the uncertainties surrounding these corrections.

Following Goff and Jordan [1988, 1989], we adopt a representation of the topographic power spectrum of the form

$$P(k, l) = 4\pi^{\nu} \frac{h_{\text{rms}}^2}{k_n k_s} \left(\frac{K^2}{k_s^2} \cos^2(\theta - \theta_s) + \frac{K^2}{k_n^2} \sin^2(\theta - \theta_s) + 1 \right)^{-(\nu+1)}, \tag{7}$$

where k_s and k_n are the wave numbers in the strike and normal directions, respectively, with $k_n \geq k_s$; θ_s measures the angle clockwise from true north to the strike direction; $\theta = \arctan(k/l)$ is the angle clockwise from true north of the wave number vector; the Hurst number ν indicates the steepness of the spectrum at high wave numbers; and h_{rms} is the root-mean-square topographic height variability, defined such that

$$\int_{-\infty}^{\infty} \int_{-\infty}^{\infty} P(k, l) dk dl = h_{rms}^2.$$

In our model experiments, the parameters (k_s, k_r, v, θ_s) are set to the constant values ($8 \times 10^{-4} \text{ rad m}^{-1}$, $3 \times 10^{-5} \text{ rad m}^{-1}$, 0.95, 20°) characteristic of abyssal hills [Goff and Arbic, 2010]. However, the sensitivity of the modeled circulation to changes in h_{rms} is investigated, as this is the most significant topographic parameter in determining G in the ocean [Scott et al., 2011].

2.2. Depth Profile of Lee Wave-Induced Mixing

A profile of the rate of turbulent kinetic energy dissipation is defined at each horizontal model grid point as

$$\epsilon(x, y, z, t) = G(x, y, t)F(z), \tag{8}$$

where F is a vertical structure function characterizing the rate of breaking of the lee waves propagating upward from the bottom. A physical model for F for lee waves, analogous to that developed for internal tides [Polzin, 2009], is not yet available, so we follow the ad hoc approach of St. Laurent et al. [2002] and prescribe F as a simple analytical function. We choose a linear vertical structure of the form

$$F(z) = \frac{2(h^* - H - z)}{h^{*2}}, \tag{9}$$

where h^* is set to 1000 m on the basis of fine and microstructure estimates of lee wave-induced ϵ in the ACC, which suggest that the dissipation is focused within 1000–2000 m of the seafloor [St. Laurent et al., 2012; Waterman et al., 2013; Sheen et al., 2013]. The vertical integral of F is 1, consistent with the complete local dissipation of lee waves (2) within a height h^* off the bottom in the same grid point, i.e., lee wave radiation and breaking are assumed to be local in the horizontal [Nikurashin and Ferrari, 2010a]. Note that this linear turbulent buoyancy flux profile implies strong near-bottom buoyancy gain, and uniform weaker compensating buoyancy loss throughout the depth range between the bottom $z = -H$ and the depth $z = -H + h^*$ where the buoyancy flux associated with the breaking internal waves is assumed to disappear (in our case the bottom 1000 m); thus the buoyancy gain from the internal wave forcing is

$$B_{iw} = -\frac{\partial B_{iw}}{\partial z} = \begin{cases} h^* B_{iw0} \delta(z+H) & \text{at the bottom,} \\ B_{iw0} & -H < z < -H+h^*, \\ 0 & z > -H+h^* \end{cases} \tag{10}$$

where

$$B_{iw0} = -B_{iw}|_{z=-H} h^{*-1} = 2G\Gamma\rho_0^{-1}h^{*-2} \tag{11}$$

is the interior diapycnal buoyancy flux convergence.

Finally, we invoke the Osborn-Cox relation [Osborn and Cox, 1972] to define a diapycnal diffusivity, κ , as

$$\kappa = \frac{\Gamma}{\rho_0} \frac{\epsilon(x, y, z, t)}{N^2(x, y, z, t)} = \frac{\Gamma}{\rho_0} \frac{G(x, y, t)F(z)}{N^2(x, y, z, t)}, \tag{12}$$

where $\Gamma = 0.2$ is a mixing efficiency for shear-driven turbulence.

3. Model Configuration and Experiments

The MITgcm [Marshall et al., 1997a, 1997b] is configured to simulate a hydrostatic flow in a three-dimensional, flat-bottom, zonally reentrant channel on a β plane, intended as an idealized representation of an ACC-like circumpolar current [similar to e.g., Abernathy et al., 2011]. The model configuration is summarized by Figure 1. The channel has a zonal extent of 1000 km, a meridional extent (L_y) of 1600 km, and a depth (H) of 3000 m. A horizontal resolution of 5 km is adopted to resolve the first baroclinic deformation radius, which is approximately 10 km in the model domain. The model has 40 vertical levels spaced at intervals ranging from 10 m near the surface to 101.5 m near the bottom. The channel is bounded by a solid wall to the south and by an open boundary to the north, where weak relaxation to an exponential temperature profile, decaying from 6°C at the surface to 0.2°C at the bottom, is applied over a 25 km-wide sponge

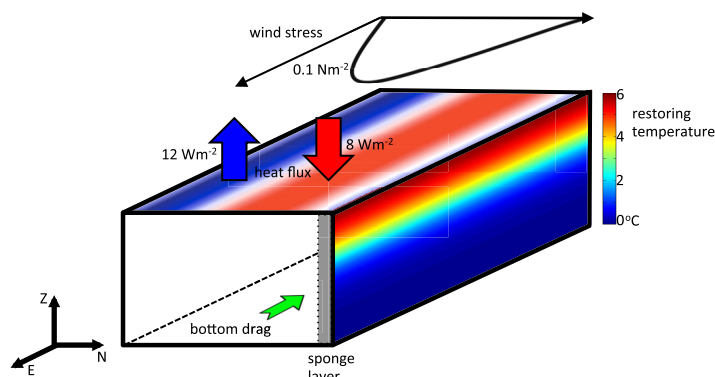


Figure 1. A schematic illustration of the model domain and the main forcing mechanisms. The unshaded areas represent solid walls.

scales facilitate a smooth forcing toward the channel interior. The model has a linear equation of state and a constant salinity of 34.5.

The model is initialized from a state of rest and a uniform temperature of 2°C, and forced with an eastward wind stress $\tau(y) = \tau_0 \sin(\pi y/L_y)$, where y is meridional distance and τ_0 is the maximum eastward wind stress. This is an idealized representation of the Southern Hemisphere westerlies. The momentum input by the wind is balanced by quadratic bottom drag with a constant coefficient of 0.01. This value, which is roughly 5 times larger than a realistic coefficient [e.g., *Naveira Garabato et al.*, 2013], prevents the development of an overly barotropic zonal flow that occurs in the absence of topographically blocked geostrophic contours [Olbers et al., 2004]. Diffusion of momentum is effected via horizontal Laplacian and biharmonic viscosity coefficients of $12 \text{ m}^2 \text{ s}^{-1}$ and $9 \times 10^8 \text{ m}^4 \text{ s}^{-1}$, respectively, and via a vertical viscosity coefficient of $3 \times 10^{-4} \text{ m}^2 \text{ s}^{-1}$. The model is also forced by a meridionally variable heat flux $Q(y) = Q_0 \cos(1.5\pi y/L_y) - \langle Q_0 \cos(1.5\pi y/L_y) \rangle$, where $Q_0 = 10 \text{ W m}^{-2}$ and angled brackets indicate a meridional average. As well as ensuring that no net buoyancy is added to the domain, this yields buoyancy loss near the channel's northern and southern edges and buoyancy gain over the center, in line with the pattern of zonal-mean buoyancy forcing of the Southern Ocean [Speer et al., 2000]. Lateral diffusion of temperature is effected by a weak skew-diffusive Gent-McWilliams [Gent and McWilliams, 1990; Griffies, 1998] flux with a diffusivity of $10 \text{ m}^2 \text{ s}^{-1}$, representing subgrid-scale eddies.

The explicit vertical diffusion is

$$D_v = \frac{\partial}{\partial z} \left(\kappa_0 \frac{\partial T}{\partial z} \right) + \frac{\partial}{\partial z} \left(\kappa_{KPP} \frac{\partial T}{\partial z} \right) + D_{KPP\text{-nonlocal}} + \frac{\partial}{\partial z} \left(\kappa_{iw} \frac{\partial T}{\partial z} \right), \quad (13)$$

where κ_0 is a constant background diapycnal diffusivity set to $10^{-5} \text{ m}^2 \text{ s}^{-1}$; D_{KPP} is the diffusivity associated with the K-Profile parameterization of mixing in the surface mixed layer and a Richardson number-dependent mixing in the interior [Large et al., 1994]; $D_{KPP\text{-nonlocal}}$ represents nonlocal KPP fluxes in the surface boundary layer; and κ_{iw} represents diffusion by internal wave breaking. Spurious numerical mixing also plays a role in models. In order to minimize this, we employ the second-order-moment scheme of Prather [1986], which produces low numerical equivalent-vertical diffusivities below $10^{-5} \text{ m}^2 \text{ s}^{-1}$ [Hill et al., 2012].

Additional parameters that constitute the model setup are the Coriolis parameter f ($1 \times 10^{-4} \text{ rad s}^{-1}$), its meridional gradient β ($1 \times 10^{-11} \text{ rad m}^{-1} \text{ s}^{-1}$) and the thermal expansion coefficient α ($2.4 \times 10^{-4} \text{ }^\circ\text{C}^{-1}$). Further details of our model setup can be obtained by downloading our model configuration, which can be obtained from the MITgcm CVS repository (see the acknowledgement section for details).

A control run is conducted with a realistic maximum wind stress, $\tau_0 = 0.1 \text{ N m}^{-2}$, and no vertical mixing by internal wave breaking, $\kappa_{iw} = 0$. The model is spun up for 200 years to achieve statistical equilibrium, and the last 10 years of this run are analyzed. Five additional experiments are presented here in which vertical mixing associated with the generation and breaking of internal lee waves is represented by κ_{iw} . This representation is introduced at the end of year 100 of the control simulation, the model run for a further 100 years, and the last 10 years considered for analysis.

layer with a tapered (northward-decreasing and downward-decreasing) time scale of 30 days at the boundary and about 960 days at the weakest point in the deep interior. This relaxation crudely represents water mass transformations north of the channel, analogous to those in the basins north of the Southern Ocean in the real world, in a much less computationally expensive manner than explicit simulation of a northern basin and the tapering time

Five additional experiments are designed to assess the effects of small-scale topography and wind stress on the modeled circulation. Three of those simulations maintain the same wind forcing as the control run, and represent lee wave-induced diapycnal mixing with three different choices of the rms topographic height, h_{rms} , the primary topographic parameter affecting G and κ_{iw} . Spatially uniform h_{rms} values of 50, 150, and 300 m are used, respectively corresponding to relatively smooth, moderate, and abrupt topography in the Southern Ocean [Scott *et al.*, 2011]. The two final simulations adopt $h_{rms} = 150$ m, and decrease or increase τ_0 by 50% (i.e., 0.05 N m^{-2}) with respect to other runs. In all perturbation experiments, the model solves equation (12) online at each grid point and time step, and the resulting κ_{iw} field both influences and is influenced by the circulation, allowing feedbacks to develop and a mutually consistent equilibrium to be attained. Note that the zonal volume transport integrated across the channel varies minimally (by less than 5%) between experiments, and that eddy kinetic energy is insensitive to the prescribed h_{rms} yet reactive to variations in wind stress (the volume integral of eddy kinetic energy in the channel decreases or increases by 31% in the experiment with stronger or weaker winds, respectively). The insensitivity of zonal volume transport and responsiveness of eddy kinetic energy to changes in wind forcing in our model are in line with the “eddy saturated” behavior exhibited by the ACC in both observations and eddy-resolving simulations [see Rintoul and Garabato, 2013, and references therein].

4. Results

The circulation in the control run is illustrated by Figure 2. A snapshot of surface temperature and horizontal velocity (Figure 2a) reveals the presence of an eastward jet flowing along the center of the channel with peak speeds of $\sim 0.5 \text{ m s}^{-1}$, roughly coincident with the maximum zonal wind stress and shedding eddies of both polarities measuring 30–100 km across. Meridional sections of the zonal-mean and time-mean potential temperature (Figure 2b) and zonal velocity (Figure 2c) indicate that the eastward jet is a permanent circulation feature, that it has an equivalent barotropic structure, and that it is in geostrophic balance with southward-shoaling isotherms. The zonal bottom velocities in the model peak at around 9 cm s^{-1} . In the ACC, bottom geostrophic flows as large as $O(10) \text{ cm s}^{-1}$ have been found [Polzin and Firing, 1997; Naveira Garabato *et al.*, 2003], which compare well to our model results. Temperature classes colder than 0°C outcrop solely in a narrow region near the channel’s southern boundary, yet dominate the deep interior, which is weakly stratified ($N \approx 3 \times 10^{-4} \text{ rad s}^{-1}$) compared to the uppermost ~ 1000 m. Note that this abyssal stratification is somewhat smaller than the typical values $\approx 10^{-3} \text{ rad s}^{-1}$ observed on the abyssal Southern Ocean [e.g., Scott *et al.*, 2011]. Zonal-mean eddy kinetic energy (Figure 2d) peaks at $\sim 0.07 \text{ m}^2 \text{ s}^{-1}$ at the surface near the channel’s axis and decays with depth by 1–2 orders of magnitude. The Eulerian-mean flow shows a classical Deacon Cell with meridional flow restricted to a northward near-surface wind-driven Ekman drift and a frictional near-bottom southward return flow. All of these properties of the modeled flow are quantitatively similar to those of ACC jets in observations [e.g., Meijers *et al.*, 2011] and eddy-resolving general circulation models [e.g., Thompson and Naveira Garabato, 2014], suggesting that our simulation is in a broadly realistic parameter regime.

The residual overturning stream function of our model simulations is calculated as the total time-averaged meridional flow below density (here reducing to temperature, as salinity is constant) surfaces, remapped to the zonal-temporal mean depth of those surfaces, following McDougall and McIntosh [2001] and in particular Nurser and Lee [2004] (NL04). Thus the time-averaged meridional transport of fluid with temperature colder than a given θ_a across a given latitude y is

$$\tilde{\psi}(y, \theta_a) = \left\langle \int \int_{(x,z): \theta(x,y,z,t) \leq \theta_a} -v \, dx \, dz \right\rangle, \tag{14}$$

(here we differ from NL04 and use the standard sign convention where the meridional stream function decreases upward for northward flow) indicating time averaging by the angled brackets $\langle \rangle$. The time-mean meridional isopycnal stream function $\tilde{\psi}(y, \theta_a)$ may be written as a function of y and z by identifying θ_a with its time-zonal-mean height

$$\tilde{z}(y, \theta_a) = \left\langle \frac{1}{L} \int_0^L z(y, \theta_a, t) \, dx \right\rangle \tag{15}$$

at each value of y , i.e., by writing

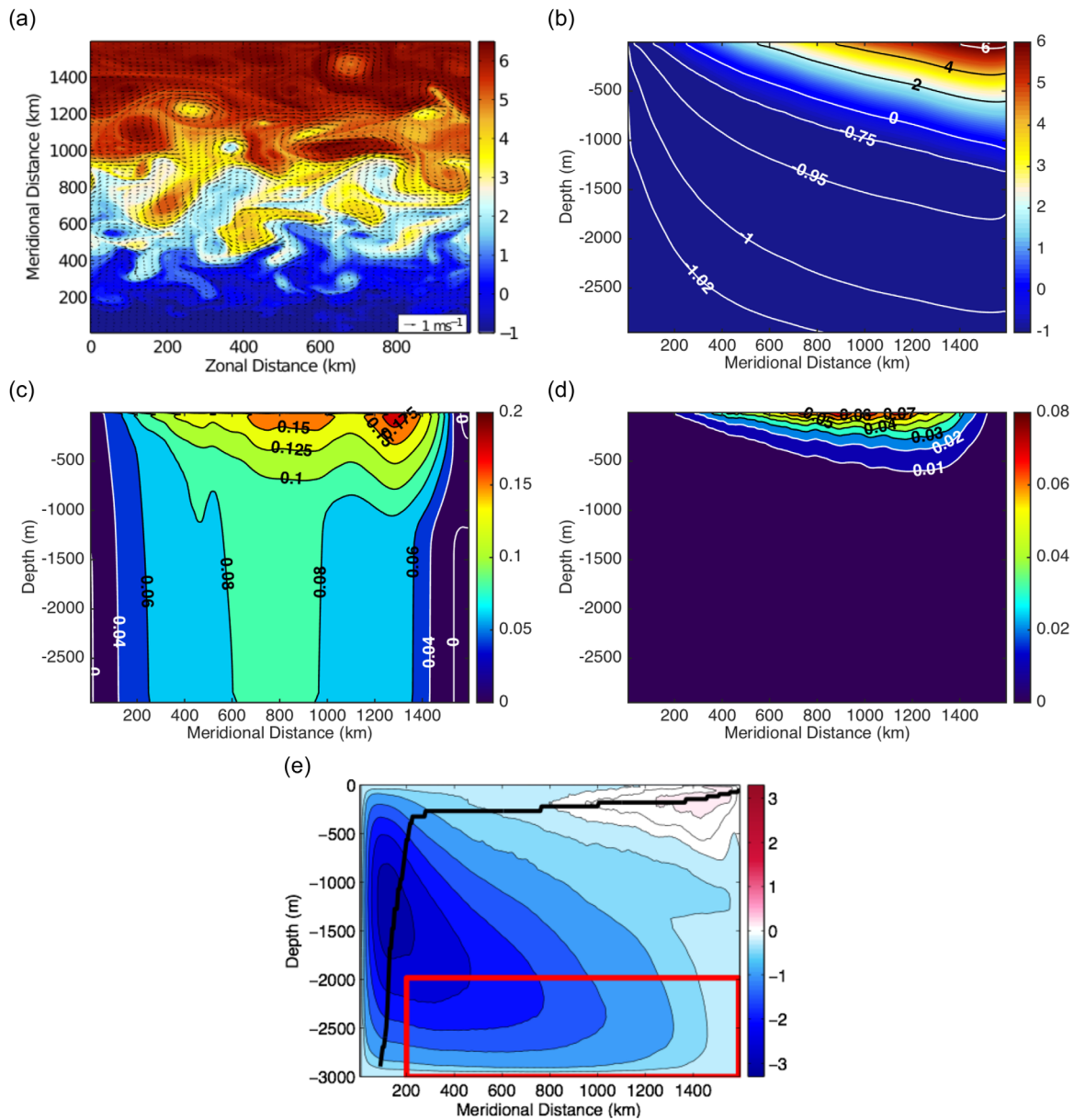


Figure 2. The control experiment: a) a snapshot and the time-mean distribution of the surface potential temperature ($^{\circ}\text{C}$, colors) overlaid with horizontal velocities (m s^{-1} , vector arrows). Average zonal-mean sections of b) the potential temperature ($^{\circ}\text{C}$), c) the zonal velocity (m s^{-1}), d) the average eddy kinetic energy ($\text{m}^2 \text{s}^{-2}$), and e) the residual overturning stream function (Sv). The thick black line in Figure 2e is the boundary of the diabatic surface forcing (DSF), as defined in the text and the red box denotes the area of averaging for the results in Figure 5.

$$\tilde{\psi}(y, \theta_a) = \tilde{\Psi}(y, \tilde{z}(y, \theta_a)), \tag{16}$$

or equivalently,

$$\tilde{\Psi}(y, z) = \tilde{\psi}(y, \tilde{\theta}(y, z)), \tag{17}$$

where $\tilde{\theta}(y, z)$ is the temperature whose zonal-temporal average height is z at latitude y . Note that this $\tilde{\theta}(y, z)$ is different to the normal Eulerian-mean $\bar{\theta}(y, z)$, most markedly near the surface where $\tilde{\theta}(y, 0)$ represents the *warmest* surface temperature that ever occurs at latitude y , not the *average* surface temperature, and similarly near the bottom $z = -H$ where $\tilde{\theta}(y, -H)$ represents the *coldest* bottom temperature that ever occurs at latitude y .

In the absence of drift (which is very small in our model runs over the averaging periods we choose), this stream function can be related to the zonally temporally averaged thickness weighted heating term

$$\tilde{Q}(y, z) = \lim_{\Delta\theta \rightarrow 0} \frac{\left\langle \iint_{(x,z): \bar{\theta} \leq \theta(x,y,z,t) \leq \bar{\theta} + \Delta\theta} \frac{D\theta}{Dt} dx dz \right\rangle}{\left\langle \iint_{(x,z): \bar{\theta} \leq \theta(x,y,z,t) \leq \bar{\theta} + \Delta\theta} dx dz \right\rangle} \quad (18)$$

according to

$$\frac{1}{L} \frac{\partial \tilde{\Psi}}{\partial y} \frac{\partial \tilde{\theta}}{\partial z} - \frac{1}{L} \frac{\partial \tilde{\Psi}}{\partial z} \frac{\partial \tilde{\theta}}{\partial y} = \tilde{Q}. \quad (19)$$

Here the y and z derivatives $\partial/\partial y$ and $\partial/\partial z$ are straightforward, taken at constant z and x , respectively. The diabatic heating terms $\frac{D\theta}{Dt}$ include surface heating and cooling, relaxation forcing, the vertical diffusion D_v from (13), numerical diffusion D_{num} , and the GM terms.

The time-averaged residual-mean circulation $\tilde{\Psi}$ in the control run exhibits a double overturning cell (Figure 2e). The channel's interior is dominated by an anticlockwise lower cell, in which ~ 3 Sv of the coldest temperature classes sink near the southern edge of the model domain, flow equatorward, and are ultimately returned upward and poleward. Note that our channel is only 1000 km long, so this transport would be very large if scaled up over the whole 18,000 km length of the ACC; this results from our strong cooling that does not occur in reality all around the ACC. A weaker (~ 0.5 Sv) clockwise upper cell is also present in the northern half of the channel. The bulk of the diathermal flow in both cells occurs in the diabatic surface layer (DSL, defined as the layer above the time-mean depth of the coldest isothermal that has entered the surface mixed layer in the period over which the circulation is temporally averaged, and indicated by the thick black line in Figure 2e), which is directly forced by the surface heat flux, and near the segment of the northern boundary sponge layer encompassed by the upper cell. Consistent with the weak interior vertical mixing prescribed in the control run, residual flow below the DSL is largely directed along isotherms. Note that the bottom (densest) residual-mean flow is northward, while the Eulerian-mean flow is everywhere southward. The structure of the simulated residual-mean circulation qualitatively resembles that of the Southern Ocean overturning, although there are significant quantitative differences between the two after the modeled circulation is scaled to the zonal extent of the Southern Ocean. Unlike in the model's case, the two Southern Ocean overturning cells are of comparable magnitude, with the upper (lower) cell being stronger (weaker) than its simulated counterpart by approximately a factor of 2 [e.g., *Garabato et al.*, 2014]. These differences are, however, of little consequence to our study, which focuses on the character and *relative* magnitude of the AABW-like lower cell's response to lee wave-induced diapycnal mixing.

We next consider the simulations with parameterized lee wave-induced diapycnal mixing. The key results for all experiments are summarized in Table 1. First, we note that the channel-averaged energy transferred to the lee wave field varies from 0.07 mW m^{-2} for the experiment with low h_{rms} (50 m), to 0.13 mW m^{-2} with moderate (150 m) and 0.18 mW m^{-2} with high (300 m) values of h_{rms} ; for the experiments with $h_{rms} = 150$ m and stronger and weaker winds, lee wave energy radiation is 0.20 and 0.04 mW m^{-2} , respectively. The increase in lee wave energy transfer with h_{rms} is much less notable than the quadratic increase ($\propto h_{rms}^2$) expected from (7); this results from an almost linear dependence of lee wave energy radiation on near-

Table 1. A Summary of the Key Results for All Experiments^a

h_{rms} (m)	Wind	Mean G (mW m^{-2})	Mean κ ($\text{m}^2 \text{ s}^{-1}$)	Mean $\tilde{\Psi}$ (Sv)	Mean θ ($^{\circ}\text{C}$)
0	Control	0	1.00×10^{-5}	-0.89	-1.00
50	Control	0.07	3.14×10^{-4}	-1.10	-0.97
150	Control	0.13	9.52×10^{-4}	-1.34	-0.96
300	Control	0.18	1.39×10^{-3}	-1.49	-0.95
150	-50%	0.04	4.57×10^{-4}	-0.76	-0.83
150	+50%	0.20	1.33×10^{-3}	-1.54	-1.03

^aThe ILW energy generation rate G is averaged over the whole channel, the mean residual overturning $\tilde{\Psi}$ and the mean potential temperature θ are averaged over the area of the red box in Figure 2e and the mean diapycnal diffusivity κ over the bottom 1000 m.

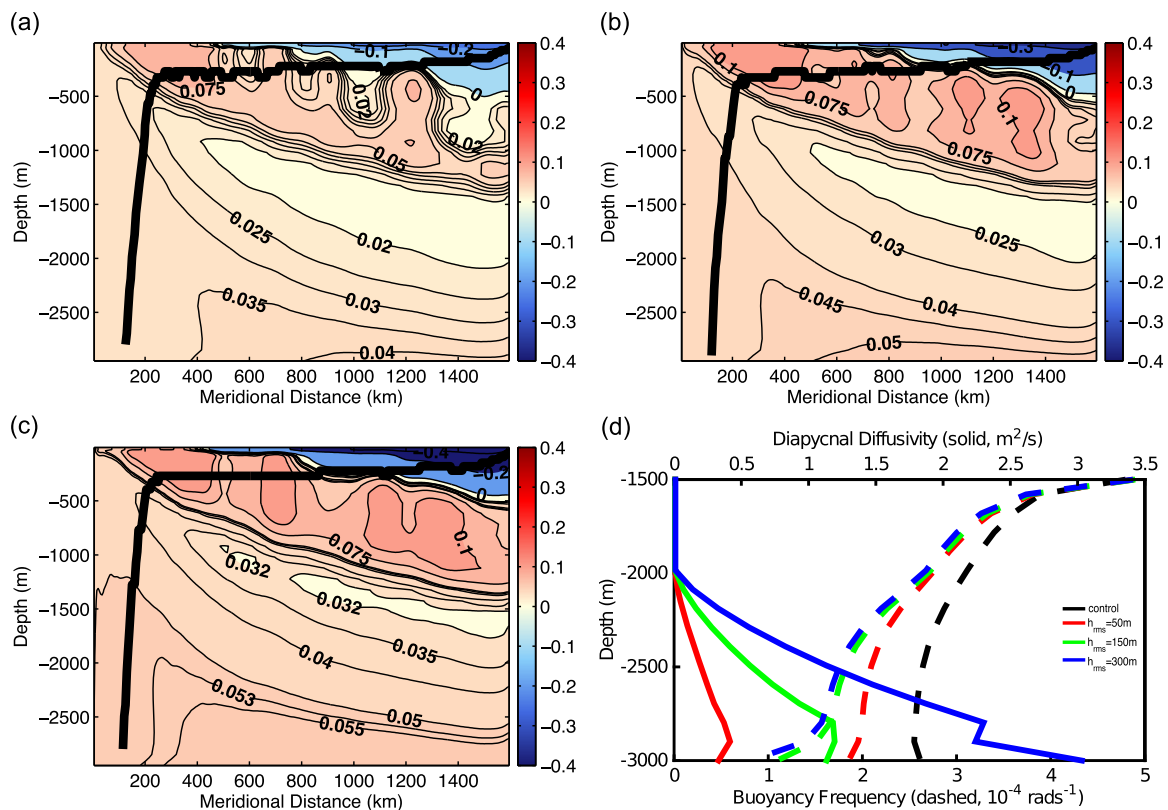


Figure 3. (a–c) The average zonal-mean section of the potential temperature difference ($^{\circ}\text{C}$) between each experiment with increasing bottom roughness parameters h_{rms} and the control experiment for $h_{rms} = 50, 150,$ and 300 m, respectively. The thick black line denotes the DSL line. (d) The channel-mean average buoyancy frequency profiles (dashed) and diapycnal diffusivity profiles (solid) for all experiments.

bottom stratification (which of course vanishes at $N=f \approx 10^{-4} \text{ rad s}^{-1}$) together with the reduction in stratification (see below) driven by the intensified near-bottom mixing. Since the bottom stratification in these runs is so weak, the IW drag coefficient given by the theory of section 2.1, is in fact an order of magnitude less than the large value of bottom drag $C_d = 0.01$ applied in our model runs.

The effect of the inclusion of lee wave-induced diapycnal mixing on the modeled diapycnal diffusivity and stratification is synthesized for the three perturbation experiments with low (50 m), moderate (150 m), and high (300 m) values of h_{rms} (Figure 3d). The horizontally and temporally averaged profiles of κ reveal an intensification of diapycnal mixing relative to the control run in all three experiments that is more pronounced with increasing h_{rms} , with mean values of $3.14 \times 10^{-4} \text{ m}^2 \text{ s}^{-1}$ ($h_{rms} = 50$ m), $9.52 \times 10^{-4} \text{ m}^2 \text{ s}^{-1}$ ($h_{rms} = 150$ m) and $1.39 \times 10^{-3} \text{ m}^2 \text{ s}^{-1}$ ($h_{rms} = 300$ m) in the bottom 1000 m (Figure 3d). The enhancement of κ has a significant impact on the modeled stratification, which is reduced as h_{rms} increases (Figure 3d). This reduction is associated with a widespread warming of the Eulerian-temporal-mean temperature in the channel interior and the southern reaches of the DSL, and with a cooling of the upper levels near the northern edge of the domain (Figures 3a–3c). The warming has two maxima: one near the bottom, where the change in κ is largest; and another in the warmest temperature classes implicated in the lower cell, which shoal from depths of 500–1400 m at the northern boundary to the surface near the southern boundary.

The elevation of diapycnal mixing and weakening of stratification with increasing h_{rms} induce a substantial response in the residual-mean circulation (Figure 4). With rougher topography, the lower overturning cell intensifies and is progressively focused closer to the bottom, as evidenced by the dipole pattern in the residual-mean stream function anomaly. Away from the DSL, the axis of the dipole is close to a depth of 2000 m, below which lee wave-induced diapycnal mixing is confined. The upper overturning cell is essentially unchanged in all perturbation experiments. Figures 4d–4f show the increased zonally averaged heating rate relative to the control for the three runs (again calculated on isotherms and remapped to their zonal-temporal-mean height), diagnosed from the residual-mean temperature equation (19); the near-

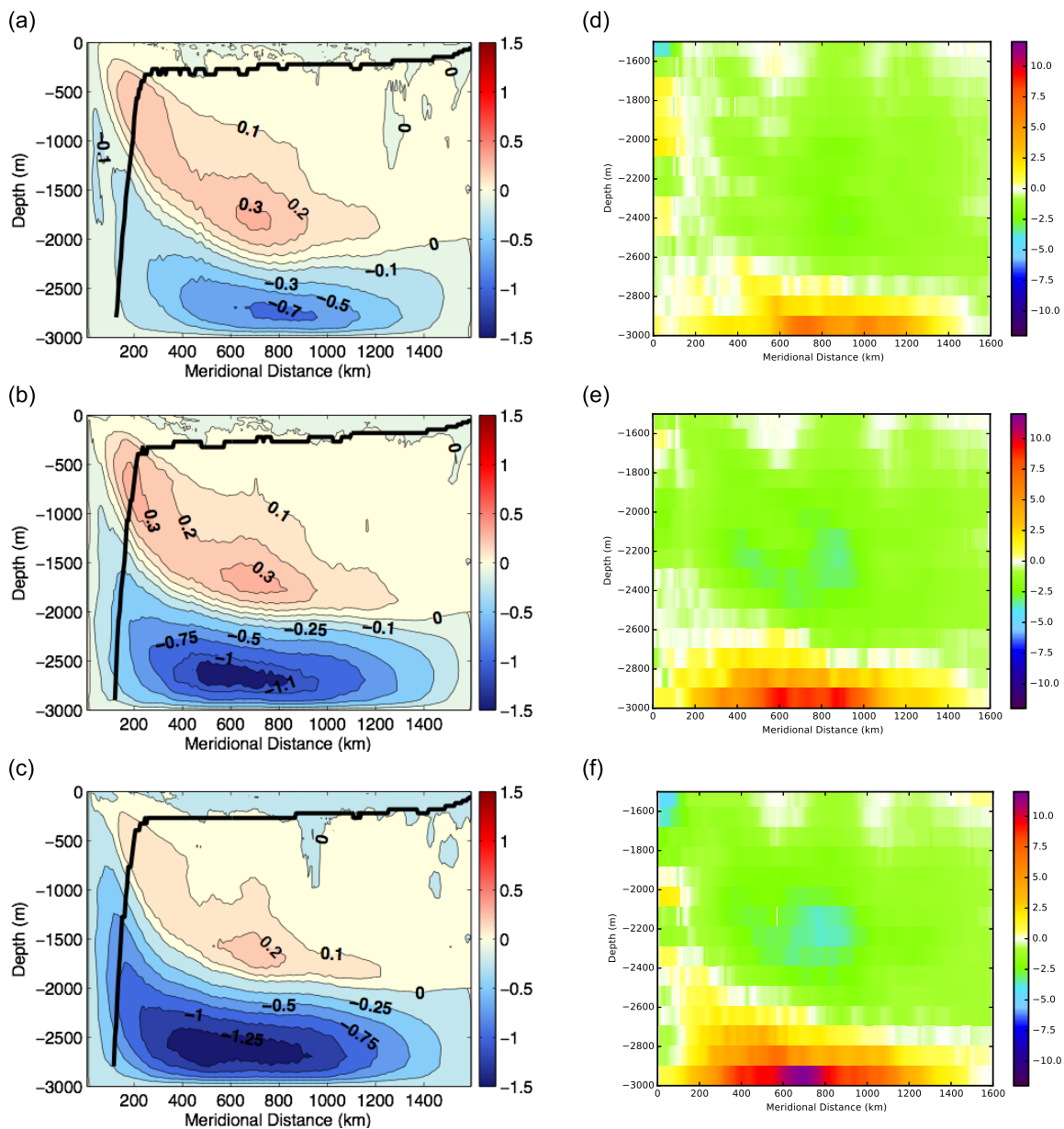


Figure 4. (a–c) The difference in the zonally integrated average residual-mean stream function (Sv) between the results from each experiment with increasing bottom roughness parameters h_{rms} and the control experiment. (a) The result for $h_{rms} = 50$ m, (b) for $h_{rms} = 150$ m, and (c) for $h_{rms} = 300$ m. The thick black line denotes the DSL line. (d–f) The differences between the residual-mean heating for the respective runs and in the control experiment in $mdeg\ ^\circ C\ yr^{-1}$. The heating is zonally averaged and thickness weighted on isotherms with given mean depth, calculated from the residual-mean stream function according to (18).

bottom warming and compensating cooling over the deepest 1000 m predicted by (10) are evident. Note that in the abyssal cell, warming happens first as the bottom waters drift northward and the water is cooled afterward as the waters drift southward and upward, so as seen in (Figures 3a–3c), the water is everywhere warmer than before despite the middepth cooling (though the temperature increase is greatest near the bottom).

An overview of how the temperature and transport of the lower overturning cell vary with h_{rms} is given by Figure 5a, in which average values of the zonal-mean temperature and the residual-mean stream function over an area spanning the bottom 1000 m of the channel away from the DSL (the red box in Figure 2e) are displayed as functions of topographic roughness. Both the temperature and strength of the lower cell increase monotonically with growing h_{rms} , although at rates that decrease with increasing roughness. This behavior is related to the relative change in κ , which increases approximately 30-fold between the control

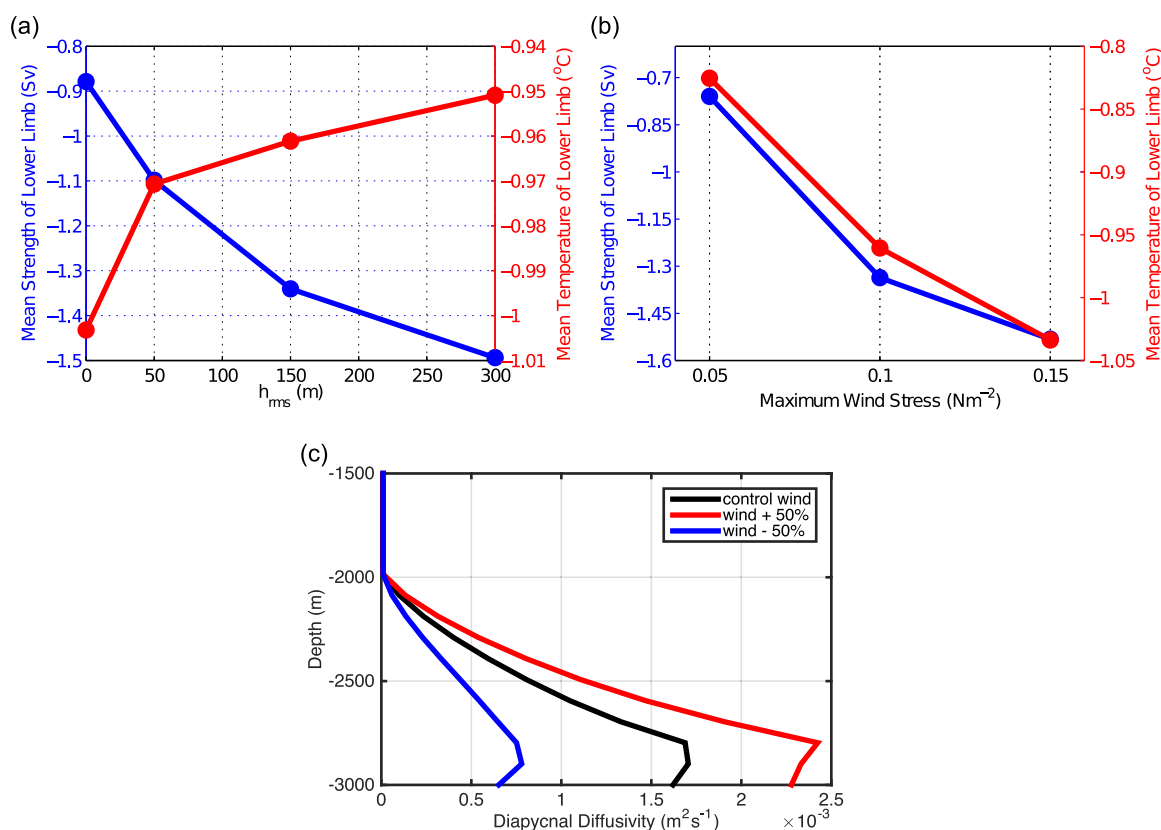


Figure 5. Summary of the evolution of the mean strength of the lower limb and the potential temperature with (a) the bottom roughness parameter for h_{rms} and (b) with the wind stress for $h_{rms} = 150$ m. The averaging was carried out over the area encompassed by the red box in Figure 2e. For $h_{rms} = 150$ m, the dependency of the diapycnal diffusivity on the wind stress is shown in Figure 5c.

simulation and the experiment with $h_{rms} = 50$ m, but only by a further factor of 1.5–3.0 for runs with higher roughness. The changes in the temperature and transport of the lower cell in the perturbation runs are significant, amounting to up to approximately 70% of the temperature difference between the downwelling and upwelling limbs of the cell and of the cell's transport in the control simulation.

The outcome of the experiments with $h_{rms} = 150$ m and perturbed eastward wind stress is synthesized by Figure 5b. A reduction (increase) of 50% in the maximum eastward wind stress leads to a substantial warming (cooling) and weakening (strengthening) of the lower overturning cell. The response is of comparable magnitude to that in the experiments with perturbed topographic roughness and is more pronounced for a reduction than for an increase in wind stress.

Finally, the dependency of the diapycnal diffusivity on the wind stress for the experiments with $h_{rms} = 150$ m is shown in Figure 5c. The response of the diffusivity to 50% weaker wind stress is approximately comparable to the result obtained by decreasing h_{rms} to 50 m, which reduces the near-bottom diffusivity by about $5 \times 10^{-4} \text{ m}^2 \text{ s}^{-1}$. Increasing the wind stress by 50% yields a less intense response in the diffusivity with an increase of about $8 \times 10^{-4} \text{ m}^2 \text{ s}^{-1}$ which is very similar to the result obtained by increasing h_{rms} to 300 m.

5. Discussion

The set of model experiments presented above indicates (i) that the transport and temperature (density) of deep meridional overturning across an ACC-like, eddy zonal flow is significantly controlled by lee wave-induced diapycnal mixing linked to the (parameterized) interaction of mesoscale eddies with small-scale topography and (ii) that this mechanism brings about a significant sensitivity of the deep residual meridional overturning to wind forcing perturbations. These findings qualitatively endorse our starting hypothesis. Next, we consider the physics underpinning the quantitative aspects of the modeled overturning behavior.

In the steady state, the variation of $\tilde{\Psi}$ with y on an isotherm of constant $\tilde{\theta}$ is given by the elegant expression, which comes immediately [Ito and Marshall, 2008] from rearranging (18) to yield

$$\left. \frac{\partial \tilde{\Psi}}{\partial y} \right|_{\tilde{\theta}} = -\tilde{Q} \left(\frac{\partial \tilde{\theta}}{\partial z} \right)^{-1}. \quad (20)$$

Ito and Marshall then integrated (20) from some value of y to the value $y_{\text{bottom}}(\tilde{\theta})$ where the isotherm strikes the bottom and the stream function $\tilde{\Psi}=0$; they assumed (i) a constant vertical diffusivity and (ii) a stratification decaying exponentially with scale depth D to give their equation (11):

$$\tilde{\Psi}(y, \tilde{\theta}) = \frac{\kappa}{D} (y - y_{\text{bottom}}(\tilde{\theta})). \quad (21)$$

This simple relation (21) describes a balance between deep-ocean buoyancy loss by advection across the isotherm and buoyancy gain by diapycnal mixing. Its direct use to diagnose our model runs would suggest much stronger abyssal cells in the runs with lee wave forcing than in the control run, given (Table 1) that the deep diffusivities over the bottom 1000 m are stronger than that in the control run by 28X in the lee waves forced with $h_{\text{rms}}=50$ m, 87X with $h_{\text{rms}}=150$ m and 126X greater with $h_{\text{rms}}=300$ m. Clearly, the increases in the overturning circulation that we see are proportionately far weaker, e.g., at 2500 m depth in the southern half of the channel, the control $\tilde{\Psi}$ is ≈ -2 Sv, increasing in strength to ≈ -2.7 Sv with $h_{\text{rms}}=50$ m, to ≈ -3.2 Sv with $h_{\text{rms}}=150$ m and to ≈ -3.4 Sv with $h_{\text{rms}}=300$ m. So why is this response so much weaker?

Most importantly (10), the downward buoyancy flux driven by the breaking lee waves attains a maximum just above the bottom boundary, and then reduces upward following the prescribed linear $F(z)$, implying an intense δ -function buoyancy (heat) gain along the bottom boundary but constant buoyancy loss (cooling) throughout the rest of the modeled deep ocean. This is consistent with a recent study by *de Lavergne et al.* [2015], who find that the bottom-most waters in the Southern Ocean gain buoyancy in the presence of near-bottom mixing whereas the overlying waters experience a buoyancy loss. This pattern of strong near-bottom warming and compensating cooling over the bottom 1000 m implies by (20) a rapid decrease of $\tilde{\Psi}$ from zero just south of the grounding latitude y_{bottom} (where the isotherm lies in the region of bottom warming), followed by a gradual return toward zero over the isotherm trajectory as it slopes upward through the bottom 1000 m—this gives very different dependence of $\tilde{\Psi}$ on y to that predicted by the simple relation (21). This pattern of decreasing and then increasing stream function along the isotherms as they curve southward and upward (e.g., Figure 2a) is consistent with the counter-clockwise perturbation cells confined to the bottom 1000 m evident in Figures 4a–4c. Near-bottom buoyancy gain drives northward flow across the sloping isotherms, balanced by southward flow across shallower isotherms driven by buoyancy loss. Because of the oppositely signed, cancellatory, contributions to the stream function, the strength of the response is far weaker than would be predicted by the simple estimate (21) for consistent buoyancy gain (warming).

However, this argument does not of itself explain why the anomalous bottom cell strength does not then scale with the increasing κ_{iw} with h_{rms} . For instance, κ_{iw} increases almost by a factor of three as h_{rms} increases from 50 to 150 m, but the anomalous $\tilde{\Psi}$ strengthens only by a factor of ~ 1.4 . And (20) suggests that it should be the diabatic heating divided by the stratification that sets $\left. \frac{\partial \tilde{\Psi}}{\partial y} \right|_{\tilde{\theta}}$; since the diabatic heating/cooling is proportional to the lee wave driven buoyancy flux, it is proportional to κ_{iw} after division by the stratification. One explanation for this is that although the vertical buoyancy (temperature) gradient weakens, the horizontal buoyancy (temperature) gradient changes little with h_{rms} . In terms of (20), this means that isotherms are almost twice as steep with $h_{\text{rms}}=150$ m (mean bottom $N \approx 1.6 \times 10^{-4} \text{ rad}^{-1}$) than with $h_{\text{rms}}=50$ m (mean bottom $N \approx 2 \times 10^{-4} \text{ rad}^{-1}$). Hence, although $\left. \frac{\partial \tilde{\Psi}}{\partial y} \right|_{\tilde{\theta}}$ increases by a factor of 3, the horizontal distance moved as the isotherm moves upward through the bottom 1000 m reduces almost by a factor of two, so $\tilde{\Psi}$ only changes by a factor of ≈ 1.7 . The importance of the horizontal buoyancy gradients is also evident from direct comparison of the contributions of the $\left. \frac{\partial \tilde{\Psi}}{\partial y} \right|_{\tilde{\theta}} \cdot \frac{\partial \tilde{\theta}}{\partial z}$ and $-\left. \frac{\partial \tilde{\Psi}}{\partial z} \right|_{\tilde{\theta}} \cdot \frac{\partial \tilde{\theta}}{\partial y}$ terms to the buoyancy (heat) budget in (19).

Indeed, no circulation would develop in our scenario of a flat-bottomed ocean with lee wave mixing without the lateral advection, as bottom waters would simply warm and convectively mix with the cooling deep

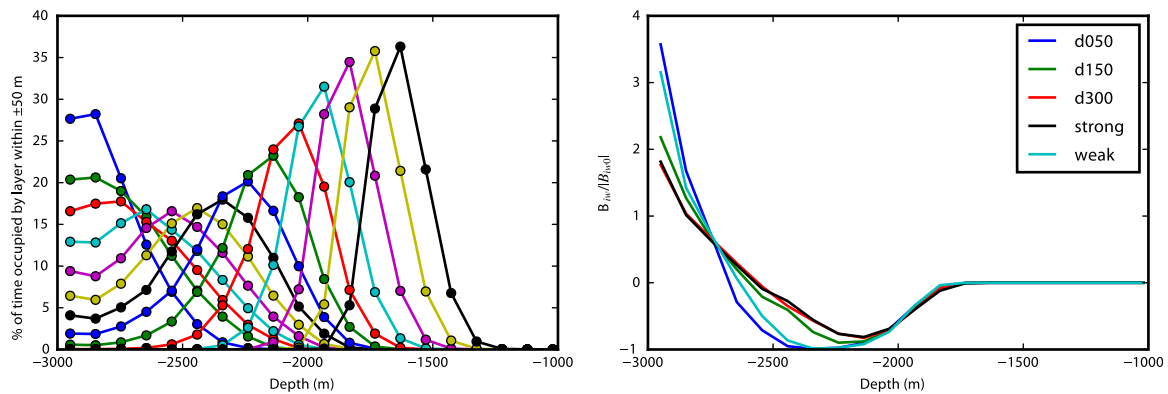


Figure 6. (a) Each colored line represents a given 100 m-thick range (specified by its color) and the dots on the line represent the percentage of time (at the middle of the channel, $y = 800$ km) that this range is occupied by isotherms with different mean depths (horizontal scale). So e.g., the leftmost (blue) line represents the % of time that the bottom 100 m is occupied by isotherms with mean depths -2950 ± 50 m, -2850 ± 50 , ... For the run with $h_{rms} = 300$ m. (b) Values of thickness weighted buoyancy input as functions of mean depth of the isotherm onto which the buoyancy input is projected (18), for the different runs, scaled by the theoretical interior buoyancy loss $|B_{iw0}|$.

waters in the bottom 1000 m. Although in reality the vertical stratification is considerably greater than in our model, our conclusion that lateral advection should make an important contribution to the density budget is, we believe, robust; isopycnal slopes in the abyssal Southern Ocean can easily reach 10^{-3} [Rintoul and Garabato, 2013], similar to the values in our model. Our findings here agree with those of Mashayek et al. [2015], who emphasize the importance of horizontal advection in the closure of the mixing driven overturning circulation.

Another process which has a major impact on the lower cell overturning driven by lee wave-induced diapycnal mixing in our simulations is the variance in the depth of the isopycnals associated with the actively eddying pseudo-ACC. Now at any one time, the lee wave-induced buoyancy flux gives a δ -function buoyancy (heat) gain near the bottom, balanced by uniform buoyancy loss B_{iw0} (cooling) throughout the deepest 1000 m. In the discretized numerical model, the deep (>1500 m depth) grid boxes have uniform thickness, 101.5 m, and the δ -function buoyancy forcing at the bottom grid level reduces to $9|B_{iw0}|$. This pattern of buoyancy forcing may be expected to drive an overturning cell with northward flow and diapycnal upwelling in a thin boundary layer adjacent to the bottom boundary, and southward flow and diapycnal downwelling elsewhere in the deepest kilometer. However, the mesoscale eddies in our simulation displace isotherms vertically, leading to a relatively wide range of isotherms being in occasional contact with the bottom boundary. Over time, this results in a spreading of the near-bottom buoyancy gain into lighter isopycnals, which appears in the residual-mean view as an “upward” spreading and broadens the time-mean vertical extent of the northward-flowing layer. This effect, in which a range of isopycnals (isotherms) come into contact with the relatively strongly forced bottom layer is directly analogous to the diabatic surface layer (DSL) where a wide range of isopycnals—all those lighter than the densest isopycnal that ever outcrops into the surface mixed layer—may be exposed to surface forcing.

To illustrate this effect of mesoscale eddies, in Figure 6a each colored line represents a given 100 m-thick range (specified by the line color), and the dots on the line represent the percentage of time that this depth range is occupied by isopycnals (isotherms) with mean depths in different 100 m ranges centered at the depths of the dots. This is along a specific zonal line, halfway across the channel (800 km north of the southern side), for the full 10 year analysis period in the simulation with $h_{rms} = 300$ m. The bottom boundary influence is found to extend upward 300–400 m, in line with the vertical scale of the northward limb of lower cell overturning (Figures 4a–4c) and diagnosed IW forcing (Figures 4d–4f). An overview of how the vertical extent of this bottom boundary influence changes across our suite of experiments, i.e., the vertical structure of the \tilde{Q} produced by the internal waves, is provided by Figure 6b. As explained above, without any isothermal movement, the scaled buoyancy input would be 9 in the bottom grid box, and -1 up to a depth of -2000 m. The upward spreading is greatest for the simulations with the roughest topography and strongest wind forcing, and most modest for the runs with the least rough topography and weakest wind forcing. This follows the weaker (stronger) stratification/more (less) energetic eddies in the former (latter)

simulations, which yields elevated (reduced) vertical displacements of isotherms. We suggest that the role of mesoscale eddies in propagating the near-bottom buoyancy gain upward over a substantial vertical scale is likely to be a realistic feature of our simulations, as it is founded on robust aspects of the modeled physics: the enhancement of diapycnal buoyancy fluxes toward the seafloor over rough topography [Waterman *et al.*, 2013; Sheen *et al.*, 2013]; the induction by eddies of vertical displacements of isopycnals of $O(100\text{ m})$ [Thompson and Naveira Garabato, 2014]; and the sensitivity of eddy kinetic energy to wind forcing [Meredith *et al.*, 2012].

6. Conclusions

The hypothesis that the eddy field's impingement on small-scale topography may regulate diapycnal mixing and meridional overturning across the deep Southern Ocean has been assessed in an idealized model. The model simulates an eddying circumpolar current coupled to a double-celled meridional overturning with properties broadly resembling those of the Southern Ocean circulation and represents lee wave-induced diapycnal mixing using an online formulation grounded on wave radiation theory. We show that the lee wave-induced diapycnal mixing generated by the simulated eddy field plays a major role in sustaining the lower overturning cell in the model, and that it underpins a significant sensitivity of this cell to wind forcing perturbations. While some simulated aspects of the overturning's quantitative dependence on diapycnal mixing are likely to be model specific, our simulations suggest that the Southern Ocean lies in a parameter regime where lee waves may be of dynamical importance to meridional overturning and point to a role of mesoscale eddies in vertically propagating the effects of enhanced near-bottom diapycnal mixing.

Our suggestion that lee wave-induced diapycnal mixing plays a significant role in sustaining deep Southern Ocean overturning resonates with the findings of several recent studies based on offline applications of wave radiation theory [Nikurashin and Ferrari, 2013; Stanley and Saenko, 2014; Mélet *et al.*, 2014]. The online formulation in our model allows us to extend these findings by unveiling how mesoscale eddies both underpin the response of deep diapycnal mixing and meridional overturning to wind forcing, and set the vertical structure of that response. Motivated by the recent and projected interdecadal intensification of the Southern Hemisphere westerlies [Thompson and Richards, 2011], many authors have investigated the eddy-mediated sensitivity of the upper cell of Southern Ocean overturning to atmospheric forcing perturbations [see e.g., Meredith *et al.*, 2012, and references therein], and suggested that flow-topography interactions are important in understanding that sensitivity [see Rintoul and Garabato, 2013, and references therein]. Our results indicate that the lower cell of Southern Ocean overturning is likely to respond to wind forcing too, via a distinct form of flow-topography interaction involving the generation and breaking of lee waves.

Like most idealized studies, ours suffers from several simplifications that may either exaggerate or understate the significance of lee wave radiation and breaking in shaping meridional overturning across the deep Southern Ocean. We have not included the effect of the IW drag on the deep flows, though we know that it is important in realistic parameter regimes [Naveira Garabato *et al.*, 2013]. Also, the circulation of newly formed bottom waters in the Southern Ocean is strongly influenced by diapycnal mixing in downslope flows and topographically guided deep boundary currents, which is likely regulated by processes distinct from lee wave radiation [e.g., Polzin *et al.*, 2014]. Further, the theory upon which our formulation of lee wave-induced diapycnal mixing is based has been suggested to significantly overestimate microstructure-derived values of κ [Waterman *et al.*, 2013; Sheen *et al.*, 2013], likely in connection with three-dimensional, finite-amplitude topographic effects [Nikurashin *et al.*, 2014]. On the other hand, the deep stratification in our model is considerably weaker than in the Southern Ocean, which significantly reduces the efficiency of energy transfer to the lee wave field. Despite these caveats, our work indicates that the Southern Ocean lies in a parameter regime where lee waves may be of dynamical importance to meridional overturning and, in revealing complex feedbacks between deep eddy flows, lee wave generation and overturning, highlights the requirement of representing lee wave processes online. Determining the degree to which our findings hold in the ocean will entail the extension of our approach to realistic general circulation models, and the precise assessment of lee wave parameterizations with observations.

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