A Toy Model for the Response of the Residual Overturning Circulation to Surface Warming

MALTE F. JANSEN

Department of the Geophysical Sciences, University of Chicago, Chicago, Illinois

LOUIS-PHILIPPE NADEAU

Institut des Sciences de la Mer de Rimouski, Université du Québec à Rimouski, Rimouski, Québec, Canada

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ABSTRACT

A simple model for the deep-ocean overturning circulation is presented and applied to study the ocean's response to a sudden surface warming. The model combines one-dimensional predictive residual advection–diffusion equations for the buoyancy in the basin and Southern Ocean surface mixed layer with diagnostic relationships for the residual overturning circulation between these regions. Despite its simplicity, the model reproduces the results from idealized general circulation model simulations and provides theoretical insights into the mechanisms that govern the response of the overturning circulation to an abrupt surface warming. Specifically, the model reproduces a rapid shoaling and weakening of the Atlantic meridional overturning circulation (AMOC) in response to surface warming, followed by a partial recovery over the following decades to centuries, and a full recovery after multiple millennia. The rapid partial recovery is associated with adjustment of the lower thermocline, which itself is shown to be accelerated by the weakened AMOC. Full equilibration instead requires adjustment of the abyssal buoyancy, which is shown to be governed by diapycnal diffusion and surface fluxes around Antarctica.

1. Introduction

Changes in the deep-ocean circulation have likely been a key player in past climatic changes, and may again be important in the future (e.g., Brovkin et al. 2007; Lund et al. 2011; Ferrari et al. 2014; Watson et al. 2015; Liu et al. 2017). The ocean's overturning circulation affects the uptake of heat and carbon during anthropogenic climate change, and can lead to potentially large changes in the regional climate via changes in the advection of warm waters (e.g., Winton et al. 2013; Frölicher et al. 2015; Marshall et al. 2015; Liu et al. 2017). Similarly, deep-ocean circulation changes have likely played an important role in the glacial cycles, by modulating deep-ocean carbon storage and hence atmospheric CO₂ concentrations (e.g., Brovkin et al. 2007; Sigman et al. 2010; Watson et al. 2015).

Our theoretical understanding of the deep-ocean circulation has improved considerably over the past decade or two, with much progress attributable to the recognition of the important role played by the Southern Ocean (e.g., Döös and Webb 1994; Toggweiler and Samuels 1995; Gnanadesikan 1999; Ito and Marshall 2008; Wolfe and Cessi 2009, 2011; Nikurashin and Vallis 2012; Mashayek et al. 2015; Thompson et al. 2016). The Southern Ocean provides a major conduit for the exchange between the surface and the deep ocean, with large amounts of deep-water upwelling along its sloping isopycnals, feeding intermediate water production to the north and Antarctic Bottom Water (AABW) production to the south (Lumpkin and Speer 2007; Marshall and Speer 2012; Talley 2013). While much of the pioneering theoretical work has focused primarily on the competing effects of wind stress and turbulent eddy transport in governing the residual overturning in the Southern Ocean, the role of surface buoyancy fluxes has more recently received increasing attention (Stewart et al. 2014; Ferrari et al. 2014; Watson et al. 2015; Sun et al. 2016; Jansen and Nadeau 2016; Jansen 2017). Specifically, it has been suggested that

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Corresponding author: Malte F. Jansen, mfi@uchicago.edu

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the buoyancy loss rate around Antarctica controls not only the abyssal stratification, but also affects the Atlantic meridional overturning circulation (AMOC), with increased buoyancy loss leading to a shoaling of North Atlantic Deep Water (NADW), which may explain inferred differences in the ocean circulation between the present and Last Glacial Maximum (LGM; Shin et al. 2003; Jansen and Nadeau 2016; Jansen 2017).

Most recent theoretical advances have been focused on steady-state equilibrium solutions. But due to the millennial time scales involved in deep-ocean adjustment, the ocean is likely not in an equilibrium state today, nor has it been in equilibrium during many of the most interesting periods in Earth's history. A better theoretical understanding of the nonequilibrium ocean response to changes in the boundary conditions on time scales from decades to millennia is thus urgently needed. The present study aims to provide a step toward such an understanding.

Using an idealized ocean general circulation model (GCM), Jansen et al. (2018) recently demonstrated the fundamental difference between the transient and equilibrium response of the ocean to surface temperature change. The equilibrium response to surface warming is found to be associated with a deepening and strengthening of the AMOC, arguably consistent with an inferred shoaling of the overturning during the cold glacial climate (e.g., Curry and Oppo 2005; Brovkin et al. 2007; Lund et al. 2011). However, the decadal to centennial response to warming is a shoaling and weakening of the AMOC, consistent with results from coupled climate simulations for the coming century (e.g., Cheng et al. 2013). While the reorganization of the equilibrium solution is driven by changes in the surface buoyancy fluxes around Antarctica, the decadal to centennial response is instead argued to be driven by warming in the NADW formation region. Idealized GCM simulations, such as those in Jansen et al. (2018), provide useful insights into the mechanisms of circulation changes on different time scales, but simplified models remain important to test our theoretical understanding and allow for wider parameter sweeps.

Various recent studies have taken steps toward the development of time-dependent theoretical models (Allison et al. 2011; Jones et al. 2011; Samelson 2011; Marshall and Zanna 2014). Allison et al. (2011), Jones et al. (2011), and Samelson (2011) build on the two-layer model of Gnanadesikan (1999) but include time tendencies to analyze the ocean's transient response to changes in the wind stress over the Southern Ocean. The models provide valuable insights into the role of circulation changes in the North Atlantic and Southern Ocean for the adjustment of the pycnocline. However, the adjustment time scales predicted by the models are considerably faster than the time scales found in

numerical simulations (e.g., Jones et al. 2011), and the two-layer model naturally cannot capture the interactions between the AMOC and the abyssal ocean, highlighted by Jansen and Nadeau (2016) and Jansen et al. (2018). A time-dependent model with much higher vertical resolution is introduced by Marshall and Zanna (2014). The model successfully reproduces the centennial-scale ocean warming seen in a coupled climate model under increased CO₂ forcing, thus providing a valuable tool to disentangle the processes that modulate ocean heat uptake in the GCM. However, the model setup employs a number of assumptions which limit its use as a predictor for ocean adjustment on time scales from decades to millennia. The buoyancy at the bottom of the ocean is fixed, which may be adequate when considering the response to warming during the first decades to centuries, but cannot account for the abyssal adjustment on longer time scales. Moreover, the isopycnal overturning circulation in the north of the basin is prescribed rather than predicted.

This manuscript introduces a new time-dependent toy model for the deep-ocean overturning circulation. The model configuration used in this study considers a single interhemispheric ocean basin, crudely representing the Atlantic, connected to a circumpolar channel in the south, representing the Southern Ocean. The evolution of the ocean's buoyancy structure is described by three one-dimensional advection–diffusion equations for the vertical buoyancy profiles in the basin and in the northern deep water formation region, as well as the meridional buoyancy structure in the Southern Ocean mixed layer. The three regions are linked using diagnostic relationships for the residual overturning circulation, following Nikurashin and Vallis (2012) and Marshall and Zanna (2014). The modular configuration of the model (which is written in Python and shared with the broader scientific community via GitHub) allows it to be easily extended or modified for future studies.

The model is applied to study the time-dependent response of the deep-ocean state and circulation to surface warming. The model accurately reproduces the idealized GCM results of Jansen et al. (2018), and provides additional insights into the mechanistic interpretation. We introduce the model's equations in section 2 and discuss the response to surface warming in section 3. Section 4 provides concluding remarks.

2. A multicolumn model for the AMOC

We will here discuss the derivation of a multicolumn toy model for the meridional overturning circulation. The present manuscript will focus on a configuration that includes a single basin with a circumpolar channel to the south. An extension to multiple basins will follow. The model's geometry is illustrated in Fig. 1. In

FIG. 1. Schematic of the single-basin multicolumn model in a depth–latitude space. Prognostic advection–diffusion equations are solved for the buoyancy profiles in the three blue shaded areas: 1) northern deep water formation region (north), 2) the basin interior (basin), and 3) surface mixed layer in the channel (SO). The predictive equations are coupled to each other using the streamfunctions computed in the green shaded areas from 4) a thermal-wind closure in the north and 5) a residual between the wind-driven and eddy-driven circulations in the channel. The dashed gray lines represent isopycnals. The latitudinal range of each region is 1) $[y_{1N}, y_{2N}]$, 2) $[y_{1B}, y_{2B}]$, and 3) $[0, L_y]$. Notice that the finite depth of the mixed layer is neglected in the computation of isopycnal slopes in the SO.

this single basin configuration, we solve predictive equations for the buoyancy profiles in three different regions: 1) the northern deep water formation region (labeled north), 2) the basin interior (labeled basin), and 3) the surface mixed layer in the model’s Southern Ocean (labeled SO). Regions 1 and 2 are represented by vertical buoyancy profiles and region 3 is represented by a meridional buoyancy profile. Each region is modeled using residual-mean advection–diffusion equations, described in sections 2a and 2d. The overturning circulation coupling the basin and the northern sinking region is computed using a thermal wind argument (box 4 in Fig. 1) and described in section 2b. The overturning transport coupling the basin and the surface mixed layer in the channel is computed using the superposition of the wind-driven and eddy-induced circulations (box 5 in Fig. 1), as discussed in section 2c. The equations are summarized in section 2e, and boundary conditions and numerical methods are discussed in sections 2f and 2g, respectively.

a. A residual-mean advection–diffusion equation for the vertical buoyancy profile

Starting with Munk (1966) the buoyancy budget in the ocean basins has often been approximated in terms of a vertical advective–diffusive balance, applied to the horizontally averaged buoyancy profile. This approximation is useful as isopycnals are typically relatively flat in the

deep ocean and away from the Southern Ocean, but it also has limitations and must break down near the bottom of the ocean (e.g., Mashayek et al. 2015). Following De Szoeke and Bennett (1993) and Young (2012) we here use an isopycnal averaging approach, which allows us to derive a residual advective–diffusive balance equation that mimics the classical vertical advective–diffusive balance albeit with some reinterpretation of the variables and coordinate system (see the appendix for a derivation):

$$\partial_t b \approx -w^\dagger \partial_z b + \partial_z (\kappa_{\text{eff}} \partial_z b) + \mathcal{B}_s. \quad (1)$$

Rather than the average buoyancy at depth z , $b(z)$ here is the inverse of $\langle z(b) \rangle$, where $\langle \rangle$ denotes the isopycnal average, that is, z is the average depth of the isopycnal $b(z)$ (Young 2012); $w^\dagger$ is the residual upwelling, κ_{eff} is an effective diapycnal diffusivity, and $\mathcal{B}_s$ represents diabatic forcing in the surface layer. The residual upwelling is given by the isopycnal volume flux convergence below the isopycnal $b(z)$ as $w^\dagger \equiv A_0^{-1} [\Psi(b(z), y_2) - \Psi(b(z), y_1)]$, with A_0 the area of the region over which we are averaging, y_1 and y_2 the latitudinal range of that region, and $\Psi(b(z), y_2)$ and $\Psi(b(z), y_1)$ the zonally integrated isopycnal overturning streamfunction at the northern and southern end of that region (we assume an average over a region that is zonally bounded by continents). In the single-basin model configuration shown in Fig. 1, Eq. (1)

is used to describe the buoyancy budget in both the basin interior (box 2 in Fig. 1, with latitudinal range $[y_{1B}, y_{2B}]$ and area $A_0 = A_B$) and the northern deep water formation region (box 1 in Fig. 1, with latitudinal range $[y_{1N}, y_{2N}]$ and area $A_0 = A_N$).

To account for isopycnals that intersect with the surface or the bottom, we define $z(b) = 0$ on the outcropped portion and $z(b) = -H$, with H the depth of the ocean, on the incropped portion of an isopycnal surface. The presence of isopycnal in- and outcrops modifies the effective diffusivity, which can be approximated as $\kappa_{\text{eff}} \approx (A_1/A_0)^2 \kappa(z)$, where A_1/A_0 is the fractional horizontal area over which the isopycnal exists (see the appendix). The factor $(A_1/A_0)^2$ accounts for the reduced area of isopycnals that intersect with a boundary, as well as for the fact that $\partial_z b$ overestimates the interior stratification, due to the contribution from isopycnal in- and outcrops. The diabatic forcing $\mathcal{B}_s$ is generally nonzero over the entire surface layer, that is, on all isopycnals that intersect with the surface. For simplicity we will here assume the surface layer over which buoyancy forcing applies to be of infinitesimal depth, which implies that we are not including the ventilated thermocline. Then $\mathcal{B}_s = 0$ in the interior and all surface fluxes are implicit in the boundary conditions. In the presence of a convectively unstable or near neutral stratification near the surface, the stratification is adjusted to a minimal value of 10^{-7} s^{-2} , which avoids numerical difficulties in the presence of zero stratification. Our results are not sensitive to the exact value chosen for the minimal stratification.

b. A thermal-wind closure for the overturning in the northern basin

To represent deep water formation in the North Atlantic, the present model will follow the approach of Nikurashin and Vallis (2012) to relate the North Atlantic Eulerian overturning transport in depth space Ψ_N^z to the latitudinal buoyancy difference between the basin interior and the deep water formation region via a thermal wind relation (box 4 in Fig. 1):

$$\partial_{zz} \Psi_N^z(z) = -f^{-1} [b_B(z) - b_N(z)]. \quad (2)$$

Here f is the Coriolis parameter, b_B is the buoyancy in the basin (box 2 in Fig. 1), and b_N is the buoyancy in the northern deep water formation region (box 1 in Fig. 1). Equation (2) more strictly represents an equation for the zonal overturning streamfunction in the north of the basin, but Nikurashin and Vallis (2012) argue that the meridional overturning transport, which occurs in the western boundary, should connect (and hence be equal) to a zonal overturning transport in the north of the basin (region 4 in Fig. 1). In practice,

Eq. (2) has been found to provide a remarkably good approximation for the meridional overturning circulation, at least when compared to idealized GCM simulations (see also Jansen et al. 2018).

Equation (2) can be solved, subject to the boundary conditions that $\Psi_N^z = 0$ at the surface and bottom of the ocean, given the buoyancy profiles in the basin $b_B(z)$ and in the NADW formation region $b_N(z)$. Nikurashin and Vallis (2012) assume b_N to be fixed at a constant value above the level where b_N matches $b_B(z)$ —the depth to which deep convection is assumed to extend—while below that depth, Ψ_N^z is assumed to be identically zero. However, Eq. (2) implies that Ψ_N^z has a turning point at the depth where $b_B - b_N = 0$, and thus $\partial_z \Psi_N^z$ (i.e., the meridional transport) has a local maximum. To allow the meridional transport to go smoothly to zero at depth, there has to be a depth range where Ψ_N^z has positive curvature and thus $b_B - b_N < 0$ at the bottom of the upper overturning cell (sketched in Fig. 1 by a negative isopycnal slope at the bottom of the lower cell in box 4). Such a reversal in the buoyancy gradient near the bottom of the overturning cell can readily be identified in idealized numerical simulations (e.g., Nikurashin and Vallis 2012; Jansen et al. 2018), although the search for a similar feature in observations is complicated by the presence of complex geometry and topography. We here explicitly solve for the buoyancy profile both in the basin and in the northern sinking region using the advective–diffusive column model described in section 2a.

The residual advective–diffusive model in section 2a requires the isopycnal overturning transport, while Eq. (2) solves for the overturning circulation in physical depth space. The isopycnal overturning circulation is computed by mapping all transport into the respective upwind buoyancy class¹ (see also Wolfe 2014):

$$\Psi_N(b) = \int_{-H}^0 \partial_z \Psi_N^z(z) \mathcal{H}[b - b_{\text{up}}(z)] dz, \quad \text{where} \\ b_{\text{up}}(z) = \begin{cases} b_N(z), & \text{if } \partial_z \Psi_N^z(z) > 0 \\ b_B(z), & \text{if } \partial_z \Psi_N^z(z) < 0. \end{cases} \quad (3)$$

c. A closure for the Southern Ocean overturning transport

The overturning circulation in the Southern Ocean (box 5 in Fig. 1) is expressed in terms of the sum of the

¹ In the numerical implementation of the model, the Heaviside step function in Eq. (3) is formulated such as to increase linearly from 0 to 1 between $b = b_{\text{up}}(z_i)$ and $b = b_{\text{up}}(z_{i+1})$, where z_i and z_{i+1} denote the levels between which $b - b_{\text{up}}(z)$ changes sign. That is, the transport between the two layers, z_i and z_{i+1} , is distributed evenly onto all buoyancy classes between $b_{\text{up}}(z_i)$ and $b_{\text{up}}(z_{i+1})$.

Ekman transport and an eddy-driven overturning circulation, following a number of previous studies (e.g., Marshall and Radko 2003; Nikurashin and Vallis 2011, 2012; Marshall and Zanna 2014):

$$\Psi_{\text{SO}} = -\frac{\tau L_x}{\rho_0 f_{\text{SO}}} + K L_x s, \quad (4)$$

where τ is the zonal wind stress, f_{SO} is the Coriolis parameter in the SO, ρ_0 is the mean ocean density, L_x is the zonal length of the Antarctic Circumpolar Current (ACC), K is an eddy “diffusivity”—or more precisely the GM (Gent and McWilliams 1990) transport coefficient—and s is the slope of the isopycnal surfaces. As in Marshall and Zanna (2014), we here use a 1D approximation to the residual overturning transport, with τ representing a characteristic wind stress over the ACC (assumed constant), and $s(b) \equiv z_B(b)/[L_y - y_{\text{SO}}(b)]$ the isopycnal slope, where $z_B(b)$ is the depth (negative) of the isopycnal in the basin, $y_{\text{SO}}(b)$ is the outcrop location of the isopycnal, and we assumed that the southern channel extends between $y = 0$ and $y = L_y$. We can compute $y_{\text{SO}}(b)$ given a monotonic surface buoyancy profile over the ACC that can be inverted to find the outcrop location.

In transient states, or if bottom water is not formed in the southern channel, the basin may have buoyancies $b_B < \min(b_{\text{SO}})$, where $\min(b_{\text{SO}})$ is the minimum surface buoyancy in the Southern Ocean channel. The slope of these non-outcropping isopycnals is set to $s = \tau/(K\rho_0 f_{\text{SO}})$, such that the residual overturning vanishes (as expected in the absence of surface water mass transformation).

Equation (4) typically yields a nonzero streamfunction value at the top and bottom of the ocean, and previous studies have employed various ways to taper the streamfunction near the boundaries (e.g., Nikurashin and Vallis 2012; Marshall and Zanna 2014). We here do not apply an explicit tapering of the streamfunction at the surface and bottom buoyancies, assuming that large transports can occur in shallow and/or weakly stratified boundary layers that cannot be represented explicitly by the model. Indeed, it can be verified that a finite transport $w^{\dagger}$ through the bottom density class in the basin (representing horizontal advection from the SO in the basin’s bottom boundary layer) is essential to obtain steady-state equilibrium solutions with vanishing diffusive buoyancy flux through the bottom boundary (see the appendix).

1) North:	$\partial_t b_N(z) = -w_N^{\dagger} \partial_z b_N + \partial_z(\kappa_{\text{eff}} \partial_z b_N)$	[cf. Eq. (1)];
2) Basin:	$\partial_t b_B(z) = -w_B^{\dagger} \partial_z b_B + \partial_z(\kappa_{\text{eff}} \partial_z b_B)$	[cf. Eq. (1)];
3) SO:	$\partial_t b_{\text{SO}}(y) = -v_{\text{SO}}^{\dagger} \partial_y b_{\text{SO}} + \partial_y(K_s \partial_y b_{\text{SO}}) + \mathcal{B}_{\text{SO}}$	[cf. Eq. (5)];
4) Thermal wind:	$\partial_{zz} \Psi_N^z = -f^{-1}(b_B - b_N)$	[cf. Eq. (2)];
5) Wind + eddy:	$\Psi_{\text{SO}} = L_x \left(-\frac{\tau}{\rho_0 f_{\text{SO}}} + K s \right)$	[cf. Eq. (4)],

d. The surface buoyancy in the channel

Equation (4) requires knowledge of the surface buoyancy profile over the SO channel, which could be prescribed as a surface boundary condition, as in Marshall and Zanna (2014). However, recent work has argued that the water mass transformation during AABW formation occurs primarily in the region of ice formation around Antarctica, where a fixed buoyancy flux condition may be more adequate (e.g., Ferrari et al. 2014; Jansen and Nadeau 2016; Jansen 2017). The idealized simulations of Jansen et al. (2018), whose results we aim to reproduce below, therefore apply a fixed buoyancy flux boundary condition in a small AABW formation region at the southern end of the domain, while applying a surface restoring boundary condition further north. To allow for such a surface flux configuration, we here introduce an explicit buoyancy budget equation for the mixed layer in the channel (box 3 in Fig. 1).

The mixed layer buoyancy budget is formulated via a horizontal residual advection–diffusion equation as

$$\partial_t b_{\text{SO}} = -v_{\text{SO}}^{\dagger} \partial_y b_{\text{SO}} + \partial_y(K_s \partial_y b_{\text{SO}}) + \mathcal{B}_{\text{SO}}, \quad (5)$$

where $v_{\text{SO}}^{\dagger} = h^{-1} L_x^{-1} \Psi_{\text{SO}}$ is the residual meridional advection, with h the mixed layer depth, K_s a horizontal mixed layer eddy diffusivity, and $\mathcal{B}_{\text{SO}}$ the surface buoyancy forcing, which is prescribed as a fixed buoyancy loss rate in a small region around the southern end of the channel, and as a surface restoring elsewhere:

$$\mathcal{B}_{\text{SO}} = \begin{cases} -h^{-1} \mathcal{F}_{\text{SO}}, & \text{for } y \leq y_{\text{fixed}} \\ v_{\text{pist}} h^{-1} (b_{\text{SO},0} - b_{\text{SO}}), & \text{for } y > y_{\text{fixed}} \end{cases}, \quad (6)$$

where $b_{\text{SO},0}$ is a prescribed latitudinal surface buoyancy profile, v_{pist} a piston velocity, and $\mathcal{F}_{\text{SO}}$ the prescribed buoyancy loss rate in the fixed flux region south of y_{fixed} . The isopycnal overturning is assumed to be constant along isopycnals in the SO, such that $\Psi_{\text{SO}}(b)$ can be mapped identically from the channel–basin interface to the surface mixed layer.

e. Summary of the model’s equations

For clarity purposes, we here summarize the model’s main equations as follows:

where $w_N^\dagger(z) = -A_N^{-1}\Psi_N$; $w_B^\dagger(z) = A_B^{-1}(\Psi_N - \Psi_{SO})$; $v_{SO}^\dagger(y) = h^{-1}\Psi_{SO}$; $s(b) \equiv z_B(b)/[L_y - y_{SO}(b)]$.

Predictive equations for the buoyancy profiles in each region b_N , b_B , and b_{SO} (boxes 1, 2, and 3 in Fig. 1, respectively) are given in the form of residual advection–diffusion equations. The different regions are coupled using diagnostic relationships for the residual streamfunctions Ψ_N and Ψ_{SO} (boxes 4 and 5 in Fig. 1, respectively), which in turn control the advective residual velocities $w_N^\dagger$, $w_B^\dagger$, and $v_{SO}^\dagger$. The slope of each isopycnal in the channel $s(b)$ is expressed in terms of the isopycnal depth in the basin region $z_B(b)$ and its outcrop location at the surface of the channel $y_{SO}(b)$.

f. Boundary conditions and parameters

The boundary conditions and parameters used for this study are chosen to mimic the idealized GCM simulations of Jansen et al. (2018). The surface buoyancy in the basin and northern deep water formation region are prescribed as $b_B(0) = 0.02 \text{ m s}^{-2}$, and $b_N(0) = -0.001 \text{ m s}^{-2}$. Notice that the absolute values of the buoyancy are irrelevant here and chosen such that the minimum restoring value in the SO is 0 m s^{-2} (see below). Assuming a thermal expansion coefficient of $\alpha = 2 \times 10^{-4} \text{ K}^{-1}$, the surface buoyancy contrast between the basin and the northern deep water formation region corresponds to about 10°C , a choice that is justified when noting that the model does not include the low-latitude ventilated thermocline.

The bottom boundary condition in the basin is set to $b_B(-H) = b_{SO}(0)$ if bottom water is entering from the SO [i.e., $\Psi_{SO}(-H + \varepsilon) < 0$, where $-H + \varepsilon$ is in practice the location of the first grid point above the bottom boundary], $b_B(-H) = b_N(-H)$ if bottom water is entering from the northern deep water formation region, and $\partial_z b_B(-H) = 0$ if no bottom water enters the abyssal basin. Similarly, we set the bottom boundary condition in the northern column such that $b_N(-H) = b_B(-H)$ if bottom water is entering from the basin, and $\partial_z b_N(-H) = 0$ if no bottom water is advected into the column.

The surface buoyancy in the channel is given by Eq. (5), where the northern boundary condition is chosen such that $b_{SO}(y = L_y) = b_B(z = 0)$. At the southern boundary, a no flux boundary condition is applied unless upwelling occurs around Antarctica [that is $\Psi_{SO}(y = \varepsilon) > 0$] in which case the boundary condition is set to the minimum buoyancy of the upwelling water, that is, $b_{SO}(y = 0) = b_B(z_0)$, where z_0 is the depth of the bottom of the upper cell at the channel–basin interface [i.e., $\Psi_{SO}(z_0) = 0$]. The width of the fixed flux region in Eq. (6) is set to $y_{\text{fixed}} = 200 \text{ km}$, and the surface restoring profile to the north is

$$b_{SO,0}(y) = 0.072\{1 - \cos[\pi(y - y_{\text{fixed}})/7.4 \times 10^6]\} \text{ m s}^{-2}. \quad (7)$$

The parameters are chosen such that the minimum buoyancy in the restoring region is $b_{SO,0}(y_{\text{fixed}}) = 0$ (i.e., just slightly larger than the buoyancy of deep water formed in the northern basin), and the surface restoring buoyancy at the northern end of the channel matches the surface buoyancy in the basin.

The diapycnal diffusivity profile, $\kappa(z)$ is based on the estimate of internal wave-driven mixing by Nikurashin and Ferrari (2013), with an added background diffusivity of $2 \times 10^{-5} \text{ m}^2 \text{ s}^{-1}$ (see Fig. 2a). The effective diffusivity κ_{eff} , which accounts for the reduced area of isopycnals in the bottom boundary layer, is obtained by tapering this profile in the bottom 500 m as $\kappa_{\text{eff}}(z) = \max[(z + H)/500, 1]^2 \kappa(z)$ whenever bottom water is advected into the respective column (basin or northern deep water formation region). Notice that the quadratic tapering profile amounts to assuming a linear decrease in the isopycnal area (see section 2a). When no bottom water is advected into the column, abyssal isopycnals are assumed to be flat, and hence $\kappa_{\text{eff}}(z) = \kappa(z)$.

Other model parameters are summarized in Table 1. Notice that the GCM simulations of Jansen et al. (2018) use a variable GM coefficient with tapering near the surface, and in addition eddy fluxes are significantly affected by resolved standing meanders, which need to be parameterized in the 1D model discussed here. As a result, K and K_s are not well constrained and here chosen to roughly mimic the GCM results in the reference simulation.

g. Numerical solution strategy

The model is solved by discretizing the three one-dimensional advection–diffusion equations and integrating numerically. A simple upwind advection scheme is employed, as numerical accuracy is not a priority, given the nature of the model. Since the vertical velocities in the basin are very small, the spurious diapycnal mixing induced by upwind differencing is negligible. Horizontal numerical diffusion in the SO mixed layer is expected to be $O(100) \text{ m}^2 \text{ s}^{-1}$, which is nonnegligible but within the uncertainty of the effective horizontal diffusivity in the GCM. We use 80 equally spaced levels ($dz = 50 \text{ m}$) for the vertical columns, and 50 meridional grid points ($dy = 40 \text{ km}$) for the surface mixed layer in the southern channel. The isopycnal mapping of the overturning circulation [Eq. (3)] uses 500 equally spaced buoyancy levels. The time step for the advection–diffusion equations is 1 month

FIG. 2. (a) Diapycnal diffusivity profile (based on Nikurashin and Ferrari 2013) used in the GCM simulations and column model (black solid), together with the effective diffusivity profile with tapering in the bottom boundary layer (blue dotted; see text for explanation). (b) Latitudinally dependent wind stress profile and (c) surface buoyancy restoring used in the GCM.

(30 days), with the diagnostic streamfunctions updated once a year.²

3. Comparison to numerical simulations

a. GCM setup

The column model discussed above will be tested against the idealized GCM simulations previously discussed in Jansen et al. (2018). The simulations employ the MITgcm (Marshall et al. 1997), in an idealized spherical domain configuration with a single 70°-wide basin extending from 48°S to 65°N, connected to a zonally re-entrant channel between 69° and 48°S, which mimics the role of the Southern Ocean. The re-entrant channel is interrupted by a sill below 3-km depth. An idealized zonally symmetric wind stress profile crudely mimics the observed large-scale zonal wind pattern (Fig. 2b). We use a linear equation of state with a thermal expansion coefficient of $\alpha = 2 \times 10^{-4} \text{K}^{-1}$ and no explicit salt, that is a single buoyancy variable effectively represents the combined density effects of temperature and salinity. While a restoring condition is applied to compute the surface buoyancy flux over most of the domain, a fixed flux is prescribed over a 2°-wide strip around Antarctica. The fixed buoyancy loss mimics the effects of brine rejection from sea ice formation around Antarctica (see also Ferrari et al. 2014; Jansen and Nadeau 2016). In Jansen et al. (2018) we show that this idealized representation of Antarctic sea ice formation allows us to adequately reproduce the results from the coupled ocean–ice

GCM. Consistent with the ocean–ice model simulations of Jansen et al. (2018), which lack sea ice in the northern basin, the buoyancy flux is only prescribed in the Southern Hemisphere. The potential role of sea ice changes in the North Atlantic for AMOC changes under global warming (Bitz et al. 2007; Zhu et al. 2015; Sévellec et al. 2017) is therefore not represented in our model. The magnitude of the prescribed area integrated buoyancy loss rate around Antarctica in the reference simulation is $5.9 \times 10^3 \text{m}^4 \text{s}^{-3}$. The surface restoring temperature profile is chosen such that the equilibrium surface buoyancy just north of the fixed flux region around Antarctica is slightly larger than the minimum buoyancy in the north of the basin (see Fig. 2c), and hence AABW formation is confined to the fixed flux region. A constant wind stress is directly prescribed at the surface using the latitudinal profile shown in Fig. 2b. The diapycnal mixing rate is horizontally homogeneous and follows the vertical profile shown in Fig. 2a (black solid). Further details about the simulations can be found in Jansen et al. (2018).

Two forcing parameters will control the transient response of the ocean circulation: 1) changes in surface temperature restoring and 2) changes in buoyancy loss around Antarctica. In addition to a “present-day-like” reference simulation we will discuss the

TABLE 1. Parameters of the column model.

$A_B = 8 \times 10^{13} \text{m}^2$	$K = 800 \text{m}^2 \text{s}^{-1}$
$A_N = 1.6 \times 10^{12} \text{m}^2$	$K_s = 400 \text{m}^2 \text{s}^{-1}$
$L_x = 4000 \text{km}$	$\tau = 0.12 \text{N m}^{-2}$
$L_y = 2000 \text{km}$	$\rho_0 = 1030 \text{kg m}^{-3}$
$h = 50 \text{m}$	$v_{\text{pist}} = 1.5 \text{m day}^{-1}$
$f = -f_{\text{SO}} = 1.2 \times 10^{-4} \text{s}^{-1}$	$\mathcal{F}_{\text{SO}} = 7.4 \times 10^{-9} \text{m}^2 \text{s}^{-3}$

² It was confirmed that updating the streamfunctions at every time step is unnecessary and would not alter our results.

FIG. 3. Buoyancy structure and meridional overturning streamfunction in the reference GCM simulation. (a) Buoyancy profile averaged between 40°S and 40°N (b_b , solid blue) and between 60° and 64°N in the northern deep water formation region (b_N , solid red), as well as the overturning streamfunction at 48°S (Ψ_{SO} , dashed cyan) and at 50°N (Ψ_N , dashed magenta). (b) Buoyancy profile at the surface of the channel, averaged over the upper 50 m, as a function of latitude (solid red), as well as the overturning streamfunction at 50-m depth (dashed cyan). (c) Depth-averaged overturning streamfunction, including parameterized eddy advection (shading; contour interval: 1 Sv) and zonally averaged buoyancy (black contours). (d) Isopycnal overturning streamfunction. The vertical coordinate has been stretched according to the average isopycnal depth in the basin. The gray vertical dashed line in (c) and (d) indicates the northern edge of the re-entrant channel region.

time-dependent response of the ocean circulation to surface warming perturbations, as well as to reduced buoyancy loss around Antarctica (mimicking reduced sea ice formation around Antarctica in a warmer climate).

b. GCM results

The overturning circulation in the reference simulation shows the familiar structure seen in the real ocean and previous idealized studies (Fig. 3c). Below shallow surface wind-driven cells, the overturning consists of two cells: an upper cell associated with northern deep water formation and an abyssal cell associated with southern bottom water formation. The buoyancy structure is qualitatively similar to the one sketched in Fig. 1, with a negative isopycnal slope in the channel region, a positive isopycnal slope in the northern deep water formation region and relatively flat isopycnals in the basin interior. Profiles of buoyancy and streamfunction corresponding roughly to the regions described in the column model are shown in Figs. 3a and 3b. The buoyancy profile in the northern sinking region (solid red in Fig. 3a) is nearly homogeneous above 2000 m, while it increases exponentially in the basin

interior (solid blue in Fig. 3a). Associated with this buoyancy difference is the upper cell in the north of the basin (dashed magenta in Fig. 3a). In the south, the streamfunction in the channel (dashed cyan in Figs. 3a and 3b) is the residual between a wind-driven and eddy-driven component, which is understood to depend on the isopycnal slope and hence the buoyancy structure both in the channel and in the basin (e.g., Marshall and Radko 2003; Nikurashin and Vallis 2012).

The main results of Jansen et al. (2018), isolating the time-dependent response of the AMOC to increased surface temperature and decreased surface buoyancy fluxes around Antarctica, are summarized in Fig. 4. In the “warming only” experiment (Fig. 4a), the surface restoring temperature is uniformly increased by 3°C (represented by a buoyancy increase of $6 \times 10^{-4} \text{ m s}^{-2}$), while the surface buoyancy loss around Antarctica is kept fixed. A rapid weakening and shoaling of the AMOC is observed over the first 10–20 years, followed by a partial recovery on a multidecadal to centennial time scale and a much slower complete recovery over several millennia. In the “ice-loss only” simulation (Fig. 4b), the surface temperature is held

FIG. 4. Hovmöller plot of the overturning circulation at 50°N for three different forcing scenarios. (a) A sudden atmospheric warming of +3°C. (b) A shutdown of surface buoyancy loss around Antarctica. (c) Combined response to warming and elimination of Antarctic buoyancy loss.

fixed while eliminating surface buoyancy forcing around Antarctica, mimicking the effect of a loss of sea ice in a large warming perturbation. In contrast to the warming-only experiment, no rapid initial response is observed, but a slow strengthening and deepening of the AMOC occurs as the stratification adjusts in the abyss. Figure 4c shows the combined effects of increased surface temperature and decreased surface buoyancy fluxes around Antarctica. The direct response to surface warming is superimposed on the slower deepening of the AMOC. This “combined warming and ice loss” experiment adequately reproduces the response of a coupled ocean–ice model to a warming perturbation (see Fig. 1 in Jansen et al. 2018).

To analyze the sensitivity of our results to the magnitude of the surface warming, an additional simulation has been performed where the prescribed surface warming is reduced by a factor of 10. Figure 5 compares the anomalous overturning circulation resulting for the large and small warming perturbations (Figs. 5a and 5b, respectively). We here focus on a homogeneous warming-only scenario, for which the equilibrium state should by construction

be identical to the initial (reference) state, except for a uniform temperature change that should not affect the dynamics (due to the use of a linear equation of state). In practice, the equilibrium circulation differs slightly from the reference case since the exact location where northern convection occurs is poorly constrained in these numerical simulations, leading to some dependence on the initial conditions. The anomalies in Fig. 5 are computed with respect to the final equilibrium state, as we focus on the disequilibrium part of the solution, but the main results are unchanged if anomalies are computed relative to the initial conditions. Despite a tenfold difference in the amplitude of the forcing, the strong and weak warming experiments share the same qualitative response. The most notable deviations from linearity include a sublinear scaling of the magnitude of the circulation response with forcing amplitude and a slightly faster recovery in the small amplitude forcing simulations, but overall, nonlinearities appear to be relatively weak.

The temporary circulation changes in the prescribed warming experiments of Fig. 5 play a crucial role in the adjustment process of the subsurface buoyancy structure.

FIG. 5. Hovmöller plot of the overturning circulation anomaly at 50°N for (a) a sudden atmospheric warming of +3°C and (b) a sudden atmospheric warming of +0.3°C. Notice the different color bars in (a) and (b).

In Jansen et al. (2018) we suggest that the lower thermocline equilibrates via changes in the overturning circulation in the northern basin, crudely following the mechanism discussed by Allison et al. (2011), Jones et al. (2011), and Samelson (2011) (hereafter collectively referred to as AJS11). The implied adjustment time scale is on the order of a few decades to a century, which is roughly consistent with the thermocline adjustment time scale in the GCM (left panel in Fig. 6). The abyssal adjustment instead is argued in Jansen et al. (2018) to be controlled by diffusion, which gives a millennial time scale, again consistent with the GCM result (right panel in Fig. 6). The adjustment time scale of both the thermocline and the abyss is again only relatively weakly dependent on the forcing amplitude. Unlike in the +3 K warming simulation, the thermocline in the +0.3 K warming simulation does not fully adjust on a centennial time scale, but this result is sensitive to the specific definition of the basin's thermocline region.

c. Column model results

In this section we first repeat the same experiments as performed with the GCM using the multicolumn model described in section 2. After confirming the model's ability to reproduce the main results from the GCM, a number of additional experiments are performed to further illuminate the mechanisms governing the ocean's adjustment.

1) REFERENCE CLIMATE AND WARMING RESPONSE

The column model qualitatively reproduces both the circulation and buoyancy structure of the reference

GCM simulation (cf. Fig. 7 to Fig. 3). The most notable difference is in the meridional structure of the residual overturning at the surface of the Southern Ocean (cf. Fig. 7b to Fig. 3b), which likely results from the simplified assumption of a constant zonal wind stress in the column model. However, the main qualitative properties—in particular a sign change at the northern edge of the fixed flux region—are reproduced. The magnitude of the upper cell is slightly weaker in the column model, transporting 9.6 Sv ($1 \text{ Sv} \equiv 10^6 \text{ m}^3 \text{ s}^{-1}$), versus 12 Sv in the GCM (at 50°N). The division between the upper and lower cell occurs around 2000–2500 m in both the GCM and the column model. The basin buoyancy profiles differ substantially in the upper ocean as the column model is missing a representation of the low-latitude ventilated thermocline, but are similar at depth. Notice that the details of the profiles shown in Fig. 3a depend on the specific latitudes chosen to compute the averages for the different regions, but the main results below the tropical thermocline are consistent between the GCM and the column model, independent of the specific choices.

The column model also reproduces the GCM's response to a sudden surface warming, and to the elimination of surface buoyancy loss around Antarctica (Fig. 8). Both the magnitude of the AMOC response and the recovery time scale are of similar order in both models. As in the GCM, the AMOC response and recovery are qualitatively similar for both small and large amplitude surface warming, although the magnitude of the overturning anomaly depends

FIG. 6. Buoyancy relative to initial conditions averaged over the basin between 40°S and 40°N. Shown are depth averages (left) over the thermocline region (above 500 m) and (right) over the abyss (below 2000 m) as a function of time. All values are normalized by the equilibrium buoyancy change. The different lines show the response to sudden surface warmings of +3 K (blue) and +0.3 K (red). Notice the different time axes in the two panels.

sublinearly on the forcing amplitude, with the non-linearity being somewhat stronger in the column model, particularly in the first couple of decades (Fig. 9).

The recovery of the thermocline and abyssal buoyancy are also broadly consistent with the GCM simulations. The most notable difference is a slight “overshoot”

with larger-than-equilibrium thermocline buoyancies after ~ 100 years, which is not observed in the GCM (Fig. 10). This overshoot, which is more pronounced in the small amplitude warming experiment, is likely associated with a persistent upper-ocean circulation anomaly associated with the delayed adjustment of the density at depth. Due to the vertically nonlocal nature of

FIG. 7. As in Fig. 3, but for the column model. Notice that the model only computes the 1D profiles shown in (a) and (b). The 2D fields in (c) and (d) are the result of an interpolation, which assumes flat or straight isopycnals in the various regions (see Fig. 1), and spatially constant up- and downwelling over the basin and in the northern deep water formation region. Notice also that the contour interval for buoyancy in (c) changes from 0.001 to 0.004 m s^{-2} at $b = 0.004 \text{ m s}^{-2}$.

FIG. 8. As in Fig. 4, but for the column model.

Eq. (2) for the overturning circulation in the basin, a reduced overturning (and hence reduced upwelling) in the upper 500 m persists even after the upper ocean buoyancy has fully adjusted. This reduced upwelling drives a temporary warming of the upper ocean beyond the prescribed surface warming. A similar overshoot is not observed in the GCM, where the upper-ocean

dynamics are significantly more complicated due to the presence of wind-driven gyres and the low-latitude thermocline. Due to the more complex upper-ocean basin dynamics, the GCM results are also sensitive to the exact depth and latitude range over which the “thermocline” average is computed, which makes a quantitative comparison with the column model difficult.

FIG. 9. As in Fig. 5, but for the column model.

FIG. 10. As in Fig. 6, but for the column model.

2) SENSITIVITY EXPERIMENTS

After confirming that the column model reproduces the main GCM results, it is instructive to probe our mechanistic understanding of the results using additional experiments, which would not readily be possible in a GCM. For this we will focus on the small amplitude warming case, as the qualitative results were found not to be sensitive to the amplitude of the forcing, and the small amplitude limit can more readily be compared to previous (linear) theories.

The goal of the sensitivity experiments is to better understand the role of dynamics in the northern basin and dynamics and thermodynamics in the SO, in the adjustment of the deep-ocean state and circulation. The scaling arguments of AJS11 highlight the importance of changes primarily in the northern overturning circulation and secondarily in the SO overturning for the adjustment of the lower thermocline, as well as the importance of SO overturning changes in the adjustment of the abyss. The scaling arguments of Jansen and Nadeau (2016) and Jansen et al. (2018) instead argue that abyssal adjustment is constrained by the balance of surface buoyancy loss around Antarctica and diffusive buoyancy gain in the basin, with the abyssal overturning playing only a passive role. To shed light on the role of these different components we will start by considering three sensitivity experiments: one where the northern overturning circulation Ψ_N^z is held fixed (in time) during the course of the experiment,³ one where the overturning circulation at the channel–basin

boundary Ψ_{SO} is held fixed,⁴ and one where the surface buoyancy in the channel b_{SO} is fixed after a uniform instantaneous increase at $t = 0$. All simulations by construction obtain the same equilibrium solution as the reference experiment, but may differ in their adjustment process.

The sensitivity experiments highlight the crucial role of changes in the northern overturning circulation in the adjustment of the lower thermocline, as well as the importance of SO thermodynamics for the adjustment of the abyss (Fig. 11). Fixing the northern overturning leads to a much slower thermocline adjustment. This result is broadly consistent with the scaling arguments of AJS11, which suggest that changes in the northern deep water formation are likely to play an important, if not dominant, role in thermocline adjustment. The dominant role of northern circulation changes is corroborated by the result that the thermocline adjustment is almost unaffected by dynamics or thermodynamics in the SO—although this result may be somewhat dependent on the model’s parameter regime, as will be elaborated below. The abyssal adjustment is largely unaffected by any changes in the overturning circulation, but is instead dependent on SO thermodynamics, with a prescribed surface buoyancy in the SO leading to a much faster adjustment of the abyss. These results are consistent with the interpretation of the abyssal adjustment process in Jansen et al. (2018). In the presence of a fixed buoyancy flux boundary condition in the outcrop region of the abyssal cell in the SO, the abyssal buoyancy

³ Notice that we are here fixing the depth-coordinate overturning Ψ_N^z , which gets mapped as before onto the buoyancy classes of the basin and northern deep water formation using the evolving buoyancy profiles in the columns. This approach ensures that at any time we are not implying any diapycnal transformations in the interior.

⁴ Notice that we fix the residual overturning circulation, as a function of depth, at the channel–basin interface. The residual transport at the surface in the SO at any time is computed by assuming an adiabatic circulation in the SO. This again avoids spurious diapycnal transformations in the interior of the SO, which would be implied if the streamfunction in the SO mixed layer was also fixed.

FIG. 11. As in Fig. 10, but comparing the response between the reference case where the overturning streamfunctions and SO surface buoyancy are fully interactive (solid red) to cases where 1) the SO overturning streamfunction is held fixed (dashed green), 2) the overturning streamfunction in the northern basin is held fixed (dotted blue), and 3) the SO surface buoyancy is held fixed after an instantaneous increase at the beginning of the simulation (dash-dotted magenta). All simulations consider the response to a surface warming of 0.3 K.

adjustment is necessarily governed by diapycnal diffusion, as changes in the SO overturning are unable to alter the buoyancy flux into the abyssal basin. A fixed surface buoyancy instead allows for a change in the advective buoyancy transport into the abyssal ocean (as implied by the scaling relations of AJS11), which leads to a reduction in the abyssal adjustment time.

Consistent with the single-basin GCM simulations, the upper overturning cell in our model is mostly diabatic, with about three quarters of the northern sinking balanced by diabatic upwelling in the basin and only about one quarter balanced by adiabatic upwelling in the Southern Ocean. However, it is worth testing whether our results hold in a more adiabatic regime, because 1) the adiabatic upwelling contribution in the real ocean is likely to be larger, primarily due to the much longer Antarctic Circumpolar Current (following standard scaling laws, the net upwelling in the channel is directly proportional to its zonal length; e.g., Marshall and Radko 2003; Nikurashin and Vallis 2012), and 2) comparison of results from the more diabatic versus adiabatic configurations provides additional insight into the role of diapycnal mixing in the adjustment process. Indeed, the results Marshall and Zanna (2014) suggest that deep-ocean heat uptake is increasingly controlled by eddy-driven advection from the SO as the ocean circulation is made more adiabatic. We analyze the adjustment process in a highly adiabatic circulation regime by reducing both the diapycnal diffusivity in the columns and the horizontal (and hence cross-isopycnal) diffusivity in the SO mixed layer by a factor of four and increasing the zonal length of the channel by a factor of three. The surface

buoyancy loss in the fixed flux region around Antarctica is chosen such that the area-integrated buoyancy loss is unchanged. The combination of reduced diapycnal mixing and increased channel length was chosen such as to obtain a roughly similar overturning in the northern basin as in our reference simulation (Fig. 12). The SO overturning instead is now much stronger, balancing most of the northern sinking, with diabatic upwelling in the basin playing a much smaller role.

At least within the context of our toy model, the major qualitative results for the adjustment of the lower thermocline and abyssal buoyancy to an instantaneous surface warming are recovered in the adiabatic setup (Fig. 13). The most notable difference is a substantial increase in the abyssal adjustment time scale, qualitatively consistent with the notion that the abyssal adjustment is controlled by diffusion. A closer inspection shows that the abyssal adjustment time scale increases only by about a factor of 2, while a simple diffusive scaling would suggest an adjustment time scale that is inversely proportional to the diapycnal diffusivity, and hence an increase by a factor of 4. The discrepancy appears to be explained by the vertical structure of the abyssal stratification that arises in the highly adiabatic limit. Following Jansen et al. (2018) the abyssal buoyancy adjusts as

$$\partial_t b'_{\text{abyss}} \sim \frac{\kappa \partial_z b'(z_{\text{int}})}{h_{\text{abyss}}}, \quad (8)$$

where b'_{abyss} is the buoyancy anomaly (relative to equilibrium) in the abyss, h_{abyss} is the vertical extent of the abyssal

FIG. 12. As in Fig. 7, but for the adiabatic AMOC configuration where the diapycnal mixing rate is reduced by a factor of 4 and the length of the SO channel is increased by a factor of 3.

cell, and $\partial_z b'(z_{\text{int}})$ is the stratification anomaly at the cell interface (see Jansen et al. 2018). The effective abyssal adjustment time scale then scales as $\tau_{\text{abyss}} \sim h_{\text{abyss}} h_{\text{grad}} / \kappa$, where h_{grad} is the depth scale associated with the buoyancy gradient at the cell interface. Since h_{grad} itself appears to decrease as κ is reduced (a sharp localized buoyancy gradient appears near the cell interface), the dependence of the adjustment time scale on κ is weaker than κ^{-1} .

The adiabatic sensitivity experiments considering fixed b_{SO} , Ψ_{SO} , and Ψ_N , respectively, also reproduce the main results found before, with two noteworthy differences. First, the sensitivity of the thermocline adjustment time scale to fixing Ψ_{SO} is increased compared to the more

diffusive configuration. The enhanced sensitivity to changes in Ψ_{SO} can be explained by the increased zonal length of the SO, which leads to a larger eddy-driven transport response for a given isopycnal slope anomaly, thus strengthening the role of SO dynamics in the adjustment process—qualitatively consistent with the arguments of AJS11 and Marshall and Zanna (2014). Second, the difference in the abyssal adjustment time scale between the full simulation and the “fixed b_{SO} ” experiment is further increased in the adiabatic limit. As discussed above, the abyssal adjustment in the fully interactive simulation (which uses a fixed-flux boundary condition around Antarctica) is controlled by diapycnal

FIG. 13. As in Fig. 11, but for adiabatic AMOC configuration.

diffusion, but it is accelerated via changes in SO eddy advection in the fixed b_{SO} experiment. While diapycnal diffusion is reduced in the adiabatic setup, eddy advection is increased as a result of the longer SO channel, which explains the increased discrepancy between the adjustment time scales of the fully interactive and fixed b_{SO} experiments. The increased discrepancy between the reference and fixed b_{SO} experiments is thus consistent with the notion that eddy advection becomes the dominant conduit for abyssal buoyancy gain in a more adiabatic ocean, if the surface buoyancy is prescribed in the AABW formation region (as in Marshall and Zanna 2014). However, if we instead prescribe the surface buoyancy *flux* in the AABW formation region, the circulation cannot affect the abyssal buoyancy budget, leading to a much slower diffusive adjustment.

4. Summary and conclusions

We introduced a modular multicolumn model for the deep-ocean overturning circulation. The model configuration used in this study consists of three 1D residual advection–diffusion equations for the evolving buoyancy structure in the basin interior, the northern deep water formation region, and the Southern Ocean. The three regions are linked using diagnostic relationships for the residual overturning circulation.

The model successfully reproduces the time-dependent response of the deep-ocean circulation to changes in the surface boundary conditions simulated by an idealized GCM. Specifically, we consider the response to a sudden change in the surface buoyancy (representing surface warming), as well as to a change in the surface buoyancy loss rate around Antarctica (representing changes in sea ice formation rate and associated brine rejection). Consistent with the GCM results, surface warming leads to a rapid shoaling and weakening of the AMOC. While the AMOC partially recovers on a centennial time scale—associated with the penetration of the warm anomaly into the thermocline in the basin—full recovery of the AMOC requires adjustment of the abyssal density, which takes multiple millennia. An elimination of Antarctic buoyancy loss leads to a shutdown of AABW formation and a gradual deepening of the AMOC, again progressing over multiple millennia.

The simple model results support the mechanistic interpretation of the circulation and buoyancy evolution suggested by Jansen et al. (2018). In particular, the model highlights the important role of AMOC changes in the adjustment of the lower thermocline, as well as the importance of thermodynamics in the Southern Ocean in governing the adjustment of the abyssal buoyancy. Specifically, if the surface boundary condition in the region of AABW formation is given by a fixed buoyancy

flux (arguably a more accurate description of surface buoyancy fluxes around Antarctica than a prescribed surface buoyancy; e.g., Jansen et al. 2018), circulation changes cannot alter the buoyancy flux into the abyssal basin, leading to a purely diffusive adjustment of the abyssal buoyancy that takes multiple millennia.

The modular configuration of the presented model (which is available via GitHub) allows it to be easily modified or extended. For example, additional ocean basins can readily be added to explore the role of inter-basin exchange, as considered by Allison (2009), Thompson et al. (2016), Jones and Cessi (2016), and Ferrari et al. (2017). Similarly, model components (such as the diagnostic relationships for the residual circulation) can be exchanged to test the sensitivity of model results to different formulations. The computational efficiency of the model (which can be integrated to equilibrium within minutes) also makes it a useful tool for classroom exercises and opens the door for studies that address ocean circulation over a wide parameter range. While certainly oversimplifying many aspects of the ocean's dynamics, we thus hope that the flexible and expandable configuration of the presented “toy model” will enable constructionist learning for students and researchers alike.

Acknowledgments. The column model code is available on Github at <https://github.com/MFJansen/PyMOC>. The MITgcm is available at <https://github.com/MITgcm/MITgcm>, and the configuration files for the simulations discussed in this study are available upon request from L.-P.N. M.F.J. acknowledges support from the NSF under Award OCE-1536454. The authors declare no conflicts of interest.

APPENDIX

Derivation of the Residual Advective–Diffusive Balance

Using the Boussinesq approximation, the continuity equation in isopycnal coordinates can be written as (e.g., Young 2012)

$$\partial_t \sigma + \nabla_h \cdot (\sigma \mathbf{u}) = -\partial_b (\sigma \mathcal{B}), \quad (\text{A1})$$

where $\sigma \equiv \partial_b z$ is the isopycnal “thickness,” ∂_t and ∇_h are the local time-derivative and horizontal nabla operator, with all derivatives taken at fixed b , $\mathbf{u} = (u, v)$ is the horizontal velocity vector,^{A1} and $\mathcal{B}$ is

^{A1} For simplicity, we here assume that isopycnal slopes are small, such that the difference between along-isopycnal and truly horizontal velocities can be neglected.

the diabatic buoyancy tendency. All variables in Eq. (A1) are taken to represent the oceanic flow field at horizontal scales $\geq O(1)$ km, with the effects of small-scale turbulence absorbed into the diabatic forcing $\mathcal{B}$. To simplify the mathematics in the presence of isopycnal outcrops, we assume Eq. (A1) to be defined at all points for $b_{\min} < b < b_{\max}$, where $b_{\min}$ and $b_{\max}$ are the minimum and maximum buoyancies over the considered domain. To extend Eq. (A1) along isopycnals that intersect with the top or bottom boundary, we simply assume those isopycnals to follow the surface or bottom of the ocean, such that $\sigma = 0$ for $b < b_{\text{bot}}(x, y)$ and $b > b_s(x, y)$, where b_s and b_{bot} are the local surface and bottom buoyancies, respectively. This approach has been introduced by Bretherton (1966) and Andrews (1983) and is also generally applied in isopycnal numerical models (e.g., Hallberg 2000).

Integrating Eq. (A1) vertically from $b_{\min}$ to some buoyancy b yields

$$\partial_t z = -\nabla_h \cdot \int_{b_{\min}}^b \sigma \mathbf{u} db' - \sigma \mathcal{B}. \quad (\text{A2})$$

The diabatic forcing can be approximated as $\sigma \mathcal{B} = -\partial_b F^b$, where

$$F^b = \begin{cases} F_s^b, & b \geq b_s(x, y) \\ -\kappa \partial_z b, & b_{\text{bot}}(x, y) < b < b_s(x, y) \\ 0, & b \leq b_{\text{bot}}(x, y) \end{cases} \quad (\text{A3})$$

is the vertical diffusive buoyancy flux, with κ representing the diapycnal turbulent eddy diffusivity; F_s^b is the surface buoyancy flux, and a no-flux boundary condition has been assumed at the bottom.

We now integrate Eq. (A2) zonally over the full extent of a basin, and meridionally between y_1 and y_2 , which yields

$$\partial_t \int_{y_1}^{y_2} \int_{x_1}^{x_2} z dx dy = \int_{x_1}^{x_2} \int_{b_{\min}}^b \sigma v db' dx \Big|_{y_2}^{y_1} + \partial_b \int_{y_1}^{y_2} \int_{x_1}^{x_2} F^b dx dy, \quad (\text{A4})$$

where x and y denote the zonal and meridional directions respectively, and a no-flux boundary condition has been assumed at the eastern and western boundaries at x_1 and x_2 . For simplicity we here assumed a Cartesian geometry and a square domain, but the results can readily be generalized to spherical coordinates and arbitrary basin geometries.

Defining the meridional isopycnal transport as $\Psi(y, b) \equiv -\int_{x_1}^{x_2} \int_{b_{\min}}^b \sigma v db' dx$, and the horizontal average as $\langle (\cdot) \rangle \equiv \int_{y_1}^{y_2} \int_{x_1}^{x_2} (\cdot) dx dy / A_0$, where $A_0 \equiv \int_{y_1}^{y_2} \int_{x_1}^{x_2} 1 dx dy$, Eq. (A4) can be written as

$$\partial_t \langle z \rangle = A_0^{-1} [\Psi(b, y_2) - \Psi(b, y_1)] + \partial_b \langle F^b \rangle. \quad (\text{A5})$$

Multiplying by $\langle \sigma \rangle^{-1}$, and noting that $\langle \sigma \rangle^{-1} \partial_t \langle z \rangle = -\partial_t b|_{\langle z \rangle}$ allows us to rewrite Eq. (A5) as a residual mean equation for the buoyancy:

$$\partial_t b|_{\langle z \rangle} = -A_0^{-1} [\Psi(b, y_2) - \Psi(b, y_1)] \partial_{\langle z \rangle} b - \partial_{\langle z \rangle} \langle F^b \rangle. \quad (\text{A6})$$

Notice that $b(\langle z \rangle)$ is not the average buoyancy at any specific depth, but rather the buoyancy of the isopycnal whose average depth is equal to $\langle z \rangle$, that is, it is the inverse function of $\langle z(b) \rangle$ (De Szeoke and Bennett 1993; Young 2012). To simplify the appearance, we define $w^\dagger \equiv A_0^{-1} [\Psi(b(\langle z \rangle), y_2) - \Psi(b(\langle z \rangle), y_1)]$, and substitute $\langle z \rangle \rightarrow z$. Equation (A6) then has the same form as the vertical advection–diffusion equation in z coordinates, albeit with some reinterpretation of the variables:

$$\partial_t b|_z = -w^\dagger \partial_z b - \partial_z \langle F^b \rangle. \quad (\text{A7})$$

To solve Eq. (A7), given the residual upwelling $w^\dagger$, we need to express $\langle F^b \rangle$ in terms of known mean-field variables. To evaluate $\langle F^b \rangle$, it makes sense to separate the horizontal integral into three components: 1) the area of the isopycnal surface that exists within the water column, 2) the area where the isopycnal is outcropped at the surface, and 3) the area where the isopycnal is incropped at the bottom (where one or two of these areas may be zero). We can then write

$$\langle F^b \rangle = -A_0^{-1} \iint_{A_1} \kappa \partial_z b dx dy + A_0^{-1} \iint_{A_2} F_s^b dx dy, \quad (\text{A8})$$

where A_1 includes the non-outcropping part of the isopycnal, where $b_{\text{bot}}(x, y) < b < b_s(x, y)$, A_2 includes the surface outcrop area, where $b > b_s(x, y)$, and no flux is imposed over the incrops, where $b < b_{\text{bot}}(x, y)$. For simplicity, we will ignore horizontal (i.e., along-isopycnal) correlations between the stratification and diapycnal diffusivity, and approximate $\kappa \approx \kappa(z)$, which yields

$$\langle F^b \rangle = -\frac{A_1}{A_0} \kappa(z) \langle \sigma^{-1} \rangle_1 + A_0^{-1} \iint_{A_2} F_s^b dx dy, \quad (\text{A9})$$

where $\langle \cdot \rangle_1$ denotes an average over the area A_1 where the isopycnal surface exists in the interior, and $A_1 \equiv \iint_{A_1} 1 dx dy$ (see Fig. A1). Assuming further that isopycnal thickness variations are relatively small over A_1 , we can approximate

$$\langle \sigma^{-1} \rangle_1 \approx \langle \sigma \rangle_1^{-1} = \left(\frac{A_0}{A_1} \langle \sigma \rangle \right)^{-1} = \frac{A_1}{A_0} \partial_z b, \quad (\text{A10})$$

FIG. A1. Example of an isopycnal intersecting the bottom. Gray shaded surfaces represent isopycnals (corresponding to the dashed gray lines in Fig. 1). Here, A_0 is the area of the entire region defined between $[x_1, x_2]$ and $[y_1, y_2]$ (green area on top). If an isopycnal out- or incrops, A_1 is the area over which the isopycnal exists in the domain (blue area at the bottom).

where the second step makes use of the fact that $\sigma = 0$ everywhere outside of A_1 .

Inserting Eqs. (A9) and (A10) into (A7) finally yields

$$\partial_t b \approx -w^\dagger \partial_z b + \partial_z (\kappa_{\text{eff}} \partial_z b) + \mathcal{B}_s, \quad (\text{A11})$$

where $\kappa_{\text{eff}} \equiv (A_1/A_0)^2 \kappa$ is an effective vertical diffusivity. The factor $(A_1/A_0)^2$ reduces the effective diffusivity on isopycnals that intersect with the surface or bottom of the ocean, which accounts for the reduced area of isopycnals that intersect with a boundary, as well as for the fact that $\partial_z b$ overestimates the interior stratification, due to the contribution from isopycnal in- and outcrops. Notice that the exact form of this prefactor depends on the assumption that the average of σ^{-1} (over the existing part of the isopycnal surface) can be approximated by the inverse of the average of σ , which may not always be a good approximation. However, the qualitative result that the effective diffusivity goes to zero as the isopycnal area vanishes, is independent of this approximation. The last term in Eq. (A11), $\mathcal{B}_s(z) \equiv -A_0^{-1} \partial_z \iint_{A_2} F_s^b dx dy$ is a diabatic forcing, which represents the effect of surface fluxes, and is generally nonzero over the surface layer, that is, on all isopycnals that outcrop somewhere in the domain.

Equation (A11) is identical in form to the well known vertical advective–diffusive balance in z coordinates, except for a reinterpretation of the coordinate system and the introduction of an effective vertical diffusivity κ_{eff} , which goes to zero at the bottom of the ocean. While seemingly a detail that only affects the equation in the bottom boundary layer, the effective diffusivity is crucial to allow for steady-state solutions with vanishing diffusive buoyancy flux through the bottom boundary. Below, we show that for $\mathcal{B}_s = 0$, Eq. (A11) has no steady-state solutions with finite stratification (i.e., $\partial_z b > 0$) in the

interior but $\partial_z b = 0$ at the bottom boundary.^{A2} Steady-state solutions with finite stratification and no flux through the bottom boundary therefore require that $\kappa_{\text{eff}} \rightarrow 0$ at the bottom of the domain.

It is worth noting that $\kappa_{\text{eff}} \rightarrow 0$ at the bottom of the domain is required but not sufficient to obtain equilibrium solutions with vanishing diffusive bottom buoyancy flux. In general, steady-state solutions to Eq. (A11) (with $\mathcal{B}_s = 0$) obey

$$\kappa_{\text{eff}}(\xi) \partial_z b(\xi) = \kappa_{\text{eff}}(\xi_0) \partial_z b(\xi_0) \exp \left(\int_{\xi_0}^{\xi} \frac{w^\dagger}{\kappa_{\text{eff}}} d\xi' \right), \quad (\text{A12})$$

where we defined $\xi = z - H$ as the distance above the bottom boundary. We now want $\kappa_{\text{eff}}(\xi_0) \partial_z b(\xi_0) \equiv F_0 > 0$ at some finite distance ξ_0 above the bottom boundary, and are interested in the buoyancy flux at the bottom, that is, the limit

$$\lim_{\xi \rightarrow 0} \kappa_{\text{eff}}(\xi) \partial_z b(\xi) = F_0 \lim_{\xi \rightarrow 0} \exp \left(\int_{\xi_0}^{\xi} \frac{w^\dagger}{\kappa_{\text{eff}}} d\xi' \right). \quad (\text{A13})$$

We find that $\lim_{\xi \rightarrow 0} \kappa_{\text{eff}}(\xi) \partial_z b(\xi) = 0$ if and only if $\lim_{\xi \rightarrow 0} \int_{\xi_0}^{\xi} (w^\dagger / \kappa_{\text{eff}}) d\xi' = -\infty$. Assuming that $w^\dagger / \kappa_{\text{eff}} \propto \xi^\alpha$ for $0 < \xi < \xi_0$, this is the case if and only if $\alpha \leq -1$. In this manuscript we use $\kappa_{\text{eff}} \propto \xi^2$ in the vicinity of the bottom boundary, while $w^\dagger$ stays finite all the way to the bottom of the (isopycnal) domain in the equilibrium solutions,^{A3} hence $\alpha = -2$ and the diffusive buoyancy flux through the bottom boundary vanishes.^{A4} Notice that a finite $w^\dagger$ at the bottom buoyancy is crucial to obtain equilibrium solutions with finite stratification and represents horizontal advection into the column.

^{A2}Notice that the formulation of the model equations in Nikurashin and Vallis (2011, 2012) suggests a vertical advective–diffusive balance equation in the basin with a no-flux bottom boundary condition. However, closer inspection of the numerical formulation shows that the model includes a horizontal advection between the basin and the last grid point of the channel. The inclusion of horizontal advection allows for the existence of steady-state solutions with finite stratification.

^{A3}Equilibrium solutions are associated with bottom water formation either in the southern channel or in the northern deep water formation column. In the former case the streamfunction, and hence $w^\dagger$, is by construction finite and increasing toward the bottom of the ocean. In the case of northern bottom water formation, the use of an upwind isopycnal advection similarly implies that bottom water enters at the minimum buoyancy.

^{A4}In practice the finite resolution of the numerical grid leads to a small amount of buoyancy loss at the bottom of the ocean, but the spurious buoyancy loss generally remains at least an order of magnitude smaller than the prescribed surface buoyancy loss rates during AABW formation.

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