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# Agulhas ring injection into the South Atlantic during glacials and interglacials

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## Abstract

Recent proxies analysis suggest that, at the end of the last glacial, there was a significant increase in the injection of Agulhas rings into the South Atlantic (SA). This brought about a dramatic increase in the salt-influx (from the Indian Ocean) into the SA helping re-start the then-collapsed meridional overturning cell (MOC), leading to the termination of the Younger Dryas (YD). Here, we propose a mechanism through which large variations in ring production take place. Using nonlinear analytical solutions for eddy shedding we show that there are restricted possibilities for ring detachment when the coast is oriented in the north-south direction. We define a critical coastline angle below which there is rings shedding and above which there is almost no shedding. In the case of the Agulhas region, the particular shape of the African continent implies that rings can be produced only when the retroflexion occurs beyond a specific latitude where the angle is critical. During glaciation, the wind stress curl (WSC) vanished at a latitude lower than that of the critical angle, which prohibited the retroflexion from producing rings. When the latitude at which the WSC vanishes migrated poleward towards its present day position, the corresponding coastline angle decreased below the critical angle and allowed for a vigorous production of rings. Simple process-oriented numerical simulations (using the Bleck and Boudra model) are in very good agreement with our results and enable us to affirm that, during the glacials, the behavior of the Agulhas Current (AC) was similar to that of the modern East Australian Current (EAC), for which the coastline slant is supercritical.

## 1 Introduction

In a recent article (Zharkov and Nof, 2007, ZN, hereafter) we examined the development of a nonlinear retroflexion eddy and constructed a solution describing the time-evolution of the ring and the mass flux going into it. Here, we shall consider the later stages in the eddies evolution – their detachment and propagation away from their gen-

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eration area. We will focus on the angle of the coastline slant, supposing the incoming current to be parallel to it, and show that, when this angle exceeds a critical value, the frequency of eddy detachment is severely restricted. We shall see that this may explain why very few Agulhas eddies were injected into the Atlantic during the Last Glaciation  
5 Maximum (LGM) and the Younger Dryas (YD),

## 1.1 Observational background

The Agulhas Current (AC) rings transport water from the Indian Ocean to the South Atlantic (SA) and, therefore, contribute to the near-surface return flow of the North Atlantic Deep Water (NADW) from the Pacific and Indian Oceans to the North Atlantic.  
10 The rings common transport represents a significant part of the meridional overturning circulation (MOC) (Gordon et al., 1987; Wejer et al., 1999; van Veldhoven, 2005; Lutjeharms, 2006) but their most important component comes from their anomalous salt content which brings in a salt anomaly five times as large as that of the Mediterranean outflow. The shedding of Agulhas rings is not a regular event. Some rings immediately  
15 split after formation, while others are recaptured by the retroflexion. Furthermore, although the typical frequency of shedding is 4–5 times a year (Schouten et al., 2002), intervals of over half-a-year without a ring-shedding event have been observed (Goni et al., 1997; van Veldhoven, 2005).

The recent paleoceanographic proxies analysis of Rau et al. (2002) and Peeters et al. (2004), suggest that the Indian-Atlantic water exchange varied greatly throughout  
20 the past 550 000 yr, having been enhanced during interglacials and strongly reduced during glacial intervals. One can tentatively suggest on the basis of Howard and Prell (1992), and Berger and Wefer (1996), that the glacial Agulhas leakage was completely shut-off due to a northward migration of the wind bands. Strictly speaking, however,  
25 this explanation is incorrect as the wind field controls the ocean interior but not the AC, i.e., rings can still propagate along the coast and penetrate into the Atlantic even when the wind stress curl vanishes north of the continental southern termination.

We note below that, even a limited reduction in ring influx could explain why the

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glacial Atlantic MOC could easily collapse. With a smaller input from Agulhas rings, the salt content of the Atlantic reduced and decreased the strength of the MOC. Because an increase in ring activity brings more salt to the Atlantic Ocean (Knorr and Lohmann, 2003), Peeters et al. (2004) suggestion that there was vigorous Agulhas ring activity at the end of each glacial might explain how a collapsed MOC restarted. The onset of increasing Agulhas leakage during late glacial conditions took place when glacial ice volume was maximal, which suggests the crucial role of Agulhas leakage in glacial terminations, timing of inter-hemispheric climate change, and the resulting resumption of the Atlantic MOC. The question that we address here is: Why were there more rings during the end of the glacial?

### 1.2 Salt balance

The importance of the Agulhas rings salt-flux to the MOC was already established by Weijer et al. (2001) and Speich et al. (2006) but it is useful to re-capture here the main aspects of the issue. The annualized volume transport associated with an average Agulhas ring is estimated to be between 0.5 and 1.5 Sv. It is not a trivial matter to estimate the salt anomaly introduced by the rings. Many conventional calculations take the salt difference between each ring and its immediate environment in the Southeastern Atlantic and calculate the contributed anomaly on this basis. This grossly underestimates the true contributed anomaly because the volume flux of the rings ( $Q_R$  which is, say, 10 Sv) is so large that the whole southeastern Atlantic is full of relatively salty water from old rings.

To correctly do the estimate, one needs to consider the salinity that the SA would have had in the absence of the rings. If one takes the SA salinity ( $S$ ) in the absence of Agulhas rings to be slightly higher than the AAIW salinity (say, 34.5 PSU) and the rings salinity to be one PSU higher, 35.5 PSU, then one finds that the salinity anomaly ( $Q_R \Delta S$ ) contributed by the rings is roughly 10 SvPSU. This is about five times the buoyancy anomaly contributed by the Mediterranean Sea or the Bering Strait.

Taking the MOC transport to be, say, 15 Sv, and applying a simple interpolation, we

find that a hypothetical removal of the entire Agulhas rings influx today (10 Sv of water 1 PSU saltier) would lower the MOC salinity by 0.7 PSU. This is about half of what would be sufficient to collapse the MOC altogether under present day conditions [see Nof and Van Gorder, 2003, (their) Fig. 2] so a simple linearization suggests that the MOC transport will be reduced by 50%.

### 1.3 The glacial-interglacial hypothesis

With basic Sverdrup dynamics, the meridional velocity within the ocean interior is proportional to the wind stress curl (WSC), whereas the zonal velocity is proportional to both the meridional gradient of the WSC and the distance from the eastern boundary of the basin. This linear Sverdrup flow occupies most of the basin, and, when the basin is closed, its net meridional transport is compensated for by a WBC flowing in the opposite direction. The meridional component of the Sverdrup transport vanishes at the latitudes where the WSC vanishes ( $\partial\tau_x/\partial y=0$ ). This implies that the flow at these latitudes is purely zonal, so that, for a closed basin, the vanishing of the WSC also defines the location of the WBC separation.

However, most ocean basins are not closed. The Indian Ocean is wide open so significant WBC-induced meridional leakages can occur across the latitude of vanishing WSC. Having said that, we should also note that, since the vanishing curl is approximately the same all around the globe, northward leakages must be compensated for by southward flow within boundary currents. Also, note that the position of zero WSC in this region is roughly at the Subtropical Convergence Zone about 45° S (de Ruijter, 1982), which is beyond the termination of the continent. In view of these, we shall take the position of vanishing WSC to be the retroflexion latitude but keep in mind that there can be leakages across it near the boundary. Furthermore, we will assume that the shift in the position of retroflexion roughly follows the shift in the WSC (Fig. 1).

The above is easier said than done, because it is difficult to determine the exact latitude of the zero WSC in the Western Indian Ocean during the LGM and the YD. On one hand, it could be inferred from Peeters et al. (2004) that the Subtropical Conver-

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gence Zone was 2–5° farther north during the YD. On the other hand, Gasse (2000) and Esper et al. (2004) argue that the shift of the westerly wind belt was much larger, as much as 25°. This is supported by the analysis of proxies from the Pacific, where the shifting process was studied via the position of the (somewhat weaker) East Australian Current (EAC). Note that, although the EAC is weaker than the AC, it is the Pacific analog of the AC. According to Martinez (1994), Kawagata (2001), Martinez et al. (2002), and Bostock et al. (2006), there was a large northward shift of the Tasman Front (a branch of the EAC) from its present latitude of 33° S to about 25° S during the last glacial. Tilburg et al. (2001) suggested that the WSC is not exactly zero at the latitude of the EAC separation, implying that the EAC does not separate completely from the coast. However, the WSC is minimal, and its gradient is maximal at the latitude of the formation of the Tasman Front, implying that the glacial/interglacial shifts of the EAC separation occurred mainly due to the shifts of the WSC.

In view of these aspects, we shall suppose that the Agulhas retroflexion latitude, during both the LGM and the YD, was between 25° S and 31° S (instead of its present position of about 38° S). We shall show that, with such a shift, the coastline slant in the neighborhood of the retroflexion during the glacials was between 50° and 70°, which severely restricted the formation of eddies. This is because the rings' long-wall drift speed was small so that they were not removed quickly enough from their formation region to avoid being recaptured. We shall see that this slower long-wall migration rate (in case of nearly meridional coastline) is due to the obvious blocking imposed by the wall.

We note here in passing that the slant of the Australian coastline at the point of the present day EAC separation is about 65°, roughly the same as that of the glacial AC. This explains why EAC rings are not usually found around the southern tip of Tasmania – the coastline slant is just too high to allow rings shedding. As in the Agulhas case, anticyclonic eddies are detached from the EAC by the pinching-off of poleward meanders. However, in contrast to the present day Agulhas case, the eddies often coalesce with the EAC again (Nilsson et al., 1977; Andrews and Scully-Power, 1976;

Nilsson and Cresswell, 1981; Sokolov and Rintoul, 2000).

The important aspect here is that none of the EAC rings (which, without the slanting coastline idea, would be the analogs of present day Agulhas rings in the Southeastern Atlantic) is found west of Tasmania.

#### 5 1.4 Present approach

We shall begin our study by examining the condition of ring shedding due to  $\beta$  (Sect. 2). We will look at the theoretical ranges of detached eddies radii, their propagation speeds, and their periods of detachment, as well as the average amount of mass flux going into the rings. To examine our glacial-interglacial hypothesis regarding the critical  
10 slant angle, we shall analyze the dependence of the above-mentioned aspects on the slant angle and other model parameters. We will also present the results obtained from numerical simulations (using a “reduced gravity” model of the Black and Boudra type), which visually elucidate two different regimes of eddy generation.

The paper is organized as follows. Section 2 is devoted to the definition of lower  
15 and upper boundaries of the ring dimensions, position and drift speed. In Sect. 3, we analyze the results of our analytical modeling, and examine the critical angles of the slant. In Sect. 4, we give the results of our numerical simulations for sub-critical and supercritical slant angles, and compare them with our analytical calculations. Finally, in Sect. 5, we summarize the results and give the conclusions regarding the glacial-  
20 interglacial hypothesis.

## 2 Shedding

### 2.1 Governing equations and the long-wall eddy propagation rate

We begin with the nonlinear momentum-flux and mass conservation equations for the retroflection eddy [also referred to as the “basic eddy” (BE)] growth as given in ZN. The

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most relevant combinations of the above are (their) Eqs. (3.21) and (3.23), which are very useful but cumbersome and, therefore, are not reproduced here. All the following formulae and computations are based on solutions of these equations for,

$$R=R(t), \Phi=\Phi(t), \quad H=H(t). \quad (1)$$

5 Here,  $R$  is the radius of developing eddy,  $\Phi$  the ratio of the mass-flux going to the BE and the total incoming mass-flux  $Q$ , and  $H$  is the thickness of the upper (moving) layer outside the retroflection area. All depend on  $Q$  (incoming mass flux),  $h_0$  (thickness of the upstream upper layer along the coast),  $g'$  (reduced gravity),  $f$  (Coriolis parameter),  $\alpha$  (twice Rossby number),  $\beta$  (meridional gradient of  $f$ ) and  $\gamma$  (slant of the coast). We  
 10 again use the tilted coordinate system  $(\xi, \eta)$  adopted by ZN. (For convenience, all variables are defined both in the text and the Appendix.)

Following Nof (2005), we consider ring shedding due to  $\beta$  which, in the open ocean, forces the eddies westward (Nof, 1983) according to,

$$C_x \approx -\left(\frac{\alpha\beta R^2}{12}\right) \frac{\alpha(2-\alpha)f^2 R^2 + 24g'H}{\alpha(2-\alpha)f^2 R^2 + 16g'H}. \quad (2)$$

15 We then suppose that, in the case of a non-zonal wall, this velocity component along the wall is simply reduced due to the geometrical blocking of the wall,

$$C_\xi = C_x \cos \gamma. \quad (3)$$

The above is not a rigorous derivation of the long-wall drift speed but it will later be confirmed with our numerical simulations. We also note that this assumption is supported  
 20 by the numerical calculations of Arruda et al. (2004), which demonstrated that no eddy detachment occurs in the case of a meridional wall ( $\gamma=90^\circ$ ). The so-called image-effect is neglected here on the ground that there is no image-effect in the limit  $H \rightarrow 0$  (Shi and Nof, 1994) and  $H$  is not very large (relative to the maximum eddy thickness) in most of our analysis. Our numerical simulation will later support this assumption even for  $H$   
 25 that is not small, probably because the rings are formed off the wall.

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## 2.2 Lower and upper boundaries for the radii of eddies and periods of their detachment

On the basis of the downstream structure (Fig. 2a), the generation period for each individual eddy is taken as,

$$t_{fI} = (2R_f + d) / |C_{\xi f}|, \quad (4)$$

where  $d$  is the distance between two consecutive eddies, and the subscript  $f$  denotes the final value. The “lower boundary” for the final eddy size ( $R_{fI}$ ) can be obtained using the “kissing condition” (i.e.  $d=0$ , Fig. 2a). In that case, (2), (3) and (4) give,

$$t_{fI} = \frac{24}{\alpha\beta R_{fI} \cos \gamma} \frac{\alpha(2-\alpha)f^2 R_{fI}^2 + 16g'H_{fI}}{\alpha(2-\alpha)f^2 R_{fI}^2 + 24g'H_{fI}}. \quad (5)$$

Equation (5) implies that  $R_{fI} = R(t_{fI})$  and  $H_{fI} = H(t_{fI})$  and will be solved numerically.

Next, we derive the intricate “upper boundary” ( $R_{fU}$ ) for the final BE size (i.e., the detachment size). For this purpose, consider the configuration shown in Fig. 2b. During the generation period, the eddy is moving along the axis  $\xi$  with the velocity  $C_{\xi}$ , which is a function of time and is defined by (2)–(3). The displacement of the BE from

its initial position during the generation period is  $\int_0^{t_{fI}} |C_{\xi}| dt$ . We note that, when this displacement equals the final diameter of the eddy (i.e.,  $2R_f$ ), then it must be detached because, at this point in time, it osculates the already generated eddy downstream (whose radius must also be  $R_f$ ).

Since the distance between the centers of the two consecutive eddies may exceed the sum of their radii, we can place the segment of the integration contour surrounding the area of the BE between the two eddies (Fig. 2b), indicating that the formation of the second eddy is now in progress as the first one has already been fully developed

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and shed. In view of this, we write the condition of “upper boundary” in the form,

$$\int_0^{t_{fu}} |C_\xi| dt = 2R_{fu}, \quad (6)$$

which is an (integral-algebraic) equation for  $R_{fu}$ .

Physically, the “upper boundary” corresponds more directly to the detachment of eddies, whereas the “lower boundary” corresponds to a condition for the formation of an eddy chain. Consequently, the eddies can detach and propagate out of the retroflection area if the condition,

$$R_{fl} \leq R_{fu} \quad (7)$$

is satisfied. This condition is certainly valid for the “imbalance paradox” (Nof, 2005), for which it is easy to show that when  $R_{fl}$  significantly exceeds  $R_i$  (the initial radius, see ZN), then  $R_{fu} = 2^{1/5} R_{fl}$ . For the upper boundary case ( $R_f = R_{fu}$ ), one obtains  $d = d_u = 2R_{fu}$ , implying that the distance between two consecutive eddies centers cannot exceed the eddy diameter.

### 3 Analysis

#### 3.1 The “lower” and “upper” boundaries and the critical angle

We (numerically) solved the corresponding nonlinear differential equations reflecting the momentum-flux and mass conservation [given as (3.21) and (3.23) in ZN], subject to conditions (1–7). For this purpose, we used:  $Q = 70 \text{ Sv}$ ,  $g' = 2 \times 10^{-3} \text{ m s}^{-2}$  and  $f = 8.8 \times 10^{-5} \text{ s}^{-1}$  (corresponding to  $35^\circ$  of latitude). We took zero and 300 m for  $h_0$ , and  $2.3 \times 10^{-11}$  and  $6 \times 10^{-11} \text{ m}^{-1} \text{ s}^{-1}$  for  $\beta$ . Alpha ( $\alpha$ ) and  $\gamma$  varied between 0.1 and 1.0, and between zero degrees and  $89^\circ$ , respectively. The results, all of which satisfy our two limiting boundaries, are shown in Fig. 3, which displays the graphs of  $R_{fl}$  and  $R_{fu}$

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as functions of  $\gamma$ . We see that, for each  $\alpha$ , the radii decrease with growing  $\gamma$ , gradually approaching each other, and finally converging. This convergence occurs when  $\gamma$  exceeds  $70^\circ$  for  $\alpha > 0.4$ , and  $60^\circ$  for  $\alpha \approx 0.2$ . When  $\alpha$  is quite small (0.1 in the figure), the curves converge very quickly and then break off. That occurs because, in this case

5 when  $\gamma$  is not very small, the  $\beta$ -force overwhelms the forces of the currents already at the initial moment (when  $R = R_i$ ) so the BE is forced into the wall instead of growing. For these conditions,  $R_{f_l}$ , and  $R_{f_u}$  cannot be defined in terms of our model. To clarify this effect further, we plotted  $R_i$ ,  $R_{f_l}$ , and  $R_{f_u}$  versus  $\alpha$  (starting from  $\alpha = 0.1$ ) for different values of  $\gamma$  (Fig. 3b). We see that the curves of  $R_{f_l}$  and  $R_{f_u}$  for  $\gamma \geq 45^\circ$  start from the

10 curve  $R_i$ , that is, they are defined for  $\alpha = 0.1$  only when  $\gamma \leq 30^\circ$ , and they are nearly convergent. This convergence extends further with increasing  $\alpha$  as the corresponding  $\gamma$  grows. In addition, although both  $R_{f_l}$  and  $R_{f_u}$  decrease with  $\alpha$ , for  $\gamma \approx 75^\circ$  this decrease becomes insignificant.

Figure 4 shows the periods of detachment  $t_{f_l}$  and  $t_{f_u}$  (left panel) and the velocities  $C_{\xi_l}$  and  $C_{\xi_u}$  of detached eddies whose radii are  $R_{f_l}$  and  $R_{f_u}$  (right panel). It is seen that the detachment period decreases with growing  $\alpha$ ; however, it increases with increasing  $\gamma$ , tending to infinity for the case of a nearly meridional wall ( $\gamma \rightarrow 90^\circ$ ). This was obviously expected because in this case  $C_\xi \rightarrow 0$ . Indeed, the curves of  $C_\xi$  converge monotonically to zero for each value of  $\alpha$  except 0.1. Convergence of curves corresponding to lower

15 and upper boundaries is seen here, too. Figures 3 and 4 suggest that eddy detachment becomes restricted when  $\gamma$  is sufficiently large, as well as when it is not large but  $\alpha$  is sufficiently small.

### 3.2 The mass flux going into the eddies

We will now estimate the ratio of the mass flux that goes into the rings and the incoming flux upstream ( $\Phi$ ). Because  $\Phi$  depends on time, it can be obtained by averaging

25 instant values over the period of eddy generation. This implies that we should operate

with two different values:

$$\Phi_l = t_{fl}^{-1} \int_0^{t_{fl}} \Phi dt, \quad \Phi_u = t_{fu}^{-1} \int_0^{t_{fu}} \Phi dt,$$

meaning that  $\Phi$  is averaged over the period of detachment.

Since  $\Phi$  grows monotonically with time only when  $\gamma=0^\circ$  and decreases for all other angles, we expect the (averaged) value for  $\gamma=0^\circ$  to be greater for longer period of formation, namely,  $\Phi_u > \Phi_l$ . Otherwise, when the wall is non-zonal, greater averaged values are expected for shorter periods, namely,  $\Phi_u < \Phi_l$ . This is clearly demonstrated in Fig. 5 showing plots of  $\Phi$  versus  $\gamma$  for different  $\alpha$  (left panel), and as functions of  $\alpha$  for different  $\gamma$  (right panel). Here, we plotted the “deadline” in analogy with that of Fig. 5 in ZN, meaning that the “vorticity paradox” occurs in the area above this line. As expected, averaged values of  $\Phi$  increase with  $\alpha$  and decrease with  $\gamma$ . If the parameter are such that the curves of  $R_{fl}$  and  $R_{fu}$  converge, then the curves of  $\Phi_u$  and  $\Phi_l$  also converge (though this convergence is weaker).

The curves of  $\Phi_u$  and  $\Phi_l$  in Fig. 5a decrease monotonically to zero, and, when considered separately for fixed  $\alpha$ , they intersect each other at the point corresponding to  $\gamma$  between  $1.5^\circ$  and  $2.5^\circ$ . The starting values are above the deadline for  $\alpha \geq 0.8$ ; nevertheless, even the maximal values, corresponding to  $\alpha=1$ , are considerably less than  $4/3$ , because the instant values of  $\Phi$  are much less than unity at the beginning of the eddy development. For  $\gamma \geq 6^\circ$ , all the curves are below the deadline. Note that, in the case of  $\alpha=0.1$ , the curves in Fig. 5a reach zero when the corresponding curves of  $R_{fl}$  and  $R_{fu}$  in Fig. 3a are terminated. In Fig. 5b,  $\Phi$  increases with  $\gamma$ . For  $\gamma=0^\circ$  (zonal wall), the corresponding waves cross the deadline for  $\alpha \approx 0.75$ , and when  $\gamma \geq 15^\circ$ , the corresponding waves are below the deadline everywhere. The rate of increase in  $\Phi$  gradually diminishes with increasing  $\gamma$ . We note here that we obtained experimental values of  $\Phi \approx 0.15$  every time  $\gamma$  exceeds  $60^\circ$ , suggesting the existence of a critical angle above which almost no eddies are detached.

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### 3.3 Varying $h_0$ and $\beta$

As mentioned, the influence of increasing  $h_0$  is not significant. When we use  $h_0=300$  m instead of zero, the decrease in  $R_{f_l}$  and  $R_{f_u}$  with  $\gamma$  becomes more significant. In the case  $\alpha>0.1$ , the corresponding curves start from almost the same values as in Fig. 3a, for  $\gamma=0^\circ$ , but tend to the range between 170 and 180 km, when  $\gamma\rightarrow 90^\circ$ . In this connection, their converging curves of  $R_{f_l}$  and  $R_{f_u}$  for  $\alpha=0.2$  intersect the ones for  $\alpha=0.4$ , exhibiting smaller limiting (for  $\gamma\rightarrow 90^\circ$ ) values. The values of  $R_{f_l}$  and  $R_{f_u}$  at the point of termination for  $\alpha=0.1$  are nearly 245 km instead of 295 km for  $h_0=0$ .

The curves of  $R_{f_l}$  and  $R_{f_u}$  versus  $\alpha$  become more spread out than in Fig. 3b, especially for  $\gamma\geq 45^\circ$  and near the curve of  $R_i$ , where they even rise slightly and reach maximal values, (240, 212, and 193 km for  $\gamma = 45^\circ, 60^\circ$ , and  $75^\circ$ , respectively). Note that  $R_i$  also decreases with increasing  $h_0$ . The effect of increasing  $h_0$  on the period of eddy generation is less significant. The increase in  $t_{f_l}$  and  $t_{f_u}$  with increasing  $\gamma$  occurs only slightly faster. Finally, the influence of  $h_0$  on the averaged values of  $\Phi$ , as well as on the velocities of detached eddies, is very weak.

As expected, changing the parameter  $\beta$  leads to quantitative changes. When we use the magnified value of  $\beta$  (i.e.,  $6\times 10^{-11} \text{ m}^{-1} \text{ s}^{-1}$ ), we find that the final eddy radius is reduced by about 20%. The boundary values of the detachment period are reduced approximately twice. The reduction in  $\Phi$  does not exceed 10%, and the velocities of the detached eddies increase by about 60%. Since the initial eddy radius does not depend on  $\beta$ , the decrease in final radii implies that there is a range of  $\alpha$  and  $\gamma$  for which the BE cannot grow at all. For example, when  $\alpha=0.1$ , the considered parameters of the detached eddies could be defined only in the case of a zonal wall ( $\gamma=0^\circ$ ). When  $\gamma=75^\circ$ , they could be defined only when  $\alpha>0.3$  (instead of  $\alpha>0.17$  for the natural value of  $\beta$ ). Therefore, the curves terminate when  $\alpha$  is less than 0.3, instead of the natural  $\beta$  termination occurring when  $\alpha=0.1$  (Fig. 3a).

Despite the reduction, the final radii are still greater than the typical observational values (see ZN, Sect. 1). For example, when  $\gamma=15^\circ$  and  $\alpha=0.15$ , then  $R_{f_l}$  and  $R_{f_u}$  are

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both approximately 275 km, but they change to 170 and 180 km, respectively, when  $\alpha$  increases to unity. At the same time,  $t_{fl}$  and  $t_{fu}$  change from 135 to 20 and 25 d, respectively. Therefore, if the potential vorticity (PV, related to  $\alpha$ , see ZN) of the eddies is nearly zero, the magnified value of  $\beta$  gives noticeably reduced values for the detachment period, which naturally is 60–90 d (Byrne et al., 1995; Schouten et al., 2000). If  $\alpha$  is about 0.2–0.3, the agreement is better.

## 4 Numerical model simulations

### 4.1 The numerical model

We used the Bleck and Boudra (1986) reduced gravity isopycnic model, the description of which was already presented in ZN and need not be repeated here. We comment here only on the simulation of the detached eddies. We initialized the retroflection from a point along the wall that was about 300 km north of the termination of the continent. However, we also conducted some experiments starting with a current retroflecting shortly after its entrance into the basin. In that case, we observed that the current behaved similarly to the case where the incoming flow was zonal (though it did so after a longer period of development). We did not obtain any plots of developing retroflection from a non-zonal wall. Our modeled time was long (about 210 d) so even when the eddies' PV was initially small, it was ultimately altered significantly by the cumulative effects of friction during the experiment. Therefore, in our quantitative comparisons, we always obtained data for nonzero PV and averaged the values of  $\alpha$  over time.

### 4.2 Varying the slant

To accelerate the detachment of the rings and make our runs more economical we chose the magnified value of  $\beta$  for most of our experiments. In the experiments with a natural  $\beta$ , the period of detachment usually exceeded 200 d but qualitative differences

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in the evolution of the BE, caused by changing  $\beta$ , were not noted. Of course, we could not perform numerical simulations in which  $\Phi$  would exceed unity.

(i) Zonal coastline: As mentioned, in the case of zonal incoming current, we were obliged to take a relatively high value for the viscosity coefficient. Although we fixed the parameters so that we can plot the thicknesses every ten days, the starting value of  $\alpha=1$  had already approached 0.4 by the time of fixing the parameters for the first plot. At that point, we identified a chain of four eddies that had already partially formed and begun to move along the wall. The first eddy was separated, and the next three were almost kissing each other (though they encountered filaments resulting from the incoming and outgoing currents). Our reduction of the friction coefficient to  $1000 \text{ m}^2\text{s}^{-1}$  resulted in a fewer kissing eddies with no filaments. We noticed almost no changes for the first four steps in time, from 10 to 40 d; and then the calculations became unstable. We conclude, therefore, that we approached the situation when  $\Phi$  was almost unity; in this case, all the incoming mass flux was contributes to eddy formation.

(ii) Sub-critical slant: The next angle,  $\gamma=15^\circ$ , also proved difficult for some numerical calculations. Because  $\Phi$  remained close to unity for  $\alpha\approx 1$ , our use of the viscosity coefficient of  $700 \text{ m}^2\text{s}^{-1}$  led to unstable calculations. Nevertheless, for this angle, the situation was better than it was for  $\gamma=0^\circ$  because, at least for  $\alpha=1$  (for which the theory suggests a short detachment period), we saw the detachment of the first two eddies. Using a larger viscosity coefficient of  $1000 \text{ m}^2\text{s}^{-1}$ , and take  $h_0=0$  after 200 d, we finally obtained the detachment of two eddies, which formed a chain downstream (Fig. 6). Before the chain formation, however, the first detached eddy (shed after about 30 d) was absorbed by a second eddy that came from behind and propagated faster. This was followed by the splitting of the merged eddies into two.

We will see later that the absorption of the first eddy occurs when  $\gamma$  is large. However, in those cases, it ultimately results in a return of the first eddy into the retroflexion area, which probably is connected with our hypothesis of restricted detachment. By contrast, for small  $\gamma$ , the occasional capturing and splitting of eddies in the model was also found in the ocean (van Veldhoven, 2005). Hence, it may either be real, or an effect of the

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large viscosity coefficient. To check this further, we prolonged the calculation so that modeled time reached 700 d. We found that the occasional capturing and splitting of eddies continued. Nevertheless, the first two eddies ultimately left the generation area (despite the increasing effect of viscosity, which led to smoothing) and a chain of five eddies was displayed.

For  $h_0=300$  m, we obtained a somewhat different situation. The detached eddies formed a chain with no capturing; however, after an intense first eddy, the second one appeared to be much weaker, and sometimes it was hardly visible. This is possibly because intense (deeper) eddies are more strongly affected by frictional forces. Similar effects can be seen for lower starting values of  $\alpha$ , but the time of eddy development is, as expected, longer. For example, when  $\alpha=0.4$  and  $h_0=0$ , the absorption of the first eddy by the second one occurs after about 200 d. When  $\alpha=0.1$ , the first eddy is not yet absorbed by day 200 but its development is almost over.

(iii) Near-critical slant: Simulations using  $\gamma=30^\circ$  and  $\gamma=45^\circ$  gave eddies that were more variable than those in the case  $\gamma=15^\circ$ . This is despite the theoretical prediction of a narrower range between the lower and upper boundaries. The first eddy was absorbed only in our simulation for  $\gamma=45^\circ$ ,  $\alpha=1$ , and  $h_0=0$ . The second one was absorbed in two simulations (with different  $h_0$ ) for  $\gamma=30^\circ$  and  $\alpha=1$ , when it was very weak. One possible explanation for the variability lies in our use of the lower viscosity coefficient of  $700 \text{ m}^2 \text{ s}^{-1}$  (which, in that case, was sufficient to ensure stability). However, a more important factor was that, despite relatively free detachment of the rings, they most often weakened gradually and, therefore, could not form a sufficiently stable chain. For a prolonged numerical simulation, the frequency of the ring capturing increased with time; as a result, after 700 d of simulation, only two rings went around the cape.

The radius of the first detached eddy increased only in our simulation for  $\gamma=30^\circ$ ,  $\alpha=0.1$ , and  $h_0=300$  m, which could be an effect of the long period of detachment. In all other cases, the radii clearly decreased, especially for  $\alpha=0.4$  and  $\alpha=1.0$ . This could be one of the main factors responsible for maintaining the distances between the rings. It is difficult to say whether this effect was a consequence of the narrowing range of the



final radius, frictional forces, or both. In any case, we conclude that the possibility of an eddy chain formation for such slants is questionable.

(iv) Super-critical slant: Simulations with a slant of  $60^\circ$ ,  $75^\circ$ , and  $90^\circ$  clearly confirmed the super-criticality because the possibility of detachment is severely limited. In some cases, such as that of  $\gamma=60^\circ$ ,  $\alpha=1$ , and  $h_0=300$  m, we did not see a detachment at all. Instead, a meander of relatively high, but variable, intensity appeared. Other simulations using  $\gamma=75^\circ$ ,  $\alpha=1$ , and  $h_0=0$ , showed almost a complete damping of the first eddy, which remained stationary after its detachment. The most typical situation is shown in Fig. 7 (for  $\gamma=60^\circ$ ,  $\alpha=1$ , and  $h_0=0$ ). After detachment, the first eddy gradually decayed before being recaptured by the meandering retroflected current behind. We note that such a situation is similar to the behavior of rings detached from the retroflection of the EAC, already described by Nilsson and Cresswell (1981), and Sokolov and Rintoul (2000). We can say, therefore, that, starting from  $\gamma=60^\circ$ , the formation of a stable chain of eddies becomes impossible. To confirm this, we again extended our calculations, so that the modeled time reached 700 d. As a result, the absorbing meander of the retroflected current was transformed into an eddy that was shed. However, later on, it was recaptured by the meander, just like the first detached eddy was. Such a process of shedding and recapturing repeated three times, and no eddies left the retroflection area. (The eddy shown in day 600 was later re-captured by day 700, but this recapturing is not displayed in Fig. 7.) In view of this, we suggest that, for such values of the slant, recapturing becomes systematic and overwhelms the formation of eddies chain.

We should comment here about the situation when  $\alpha$  is small. Our simulations with  $\alpha=0.1$  and  $\gamma=60^\circ$  and  $75^\circ$  qualitatively confirmed our theoretical prediction that the eddy cannot grow because the  $\beta$ -induced force initially exceeds the combined force of the currents (i.e., their long-shore momentum flux). In such cases, the plots showed the formation of a meander that was gradually forced into the wall. Its longitudinal dimension, along the wall, became much greater than the transverse one, negating our assumption of a nearly circular form. Eventually, an eddy of a smaller radius was

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formed from this meander. Further, the formed eddy continued decaying and exhibited almost no movement. We expected that it would finally be reabsorbed, but we did not reach that state.

### 4.3 Simulation-theory comparison

5 In addition to a qualitative analysis of ring behavior, we carried out a detailed comparison of the theoretical modeling and the numerical simulations. (For this purpose, we used the magnified value of  $\beta$ .) The parameters that we compared are the eddy radius at the moment of detachment,  $R_f$ , the eddy propagation velocity ( $C_\xi$ ), and the ratio of the mass flux going into the eddies and the incoming flux ( $\Phi$ ).

10 Two introductory comments should be made here. First, Nof and Pichevin (2001, NP, hereafter) carried out the quantitative comparison of the ratio  $q/Q$  as a function of time, where  $q$  is outgoing mass flux. In light of this, it was very important to take into account the time evolution of the parameter  $\alpha$ , which was dramatically altered by the viscosity in the numerics. For us, it was inconvenient to conduct similar analysis because  $\Phi$  was also variable in time, and we were operating mainly with its time-averaged values. Therefore, we simplified our analysis by assuming that the value of  $\Phi$ , averaged over  
15 the time of the numerical experiment, corresponded to the value of  $\alpha$  averaged in the same way, and so did with the other parameters that we considered. Second, we also averaged the value of  $C_\xi$ . Wherever possible, we used an averaging period between  
20 the moment of detachment and the last step in our experiment. However, in most simulations with supercritical values of  $\gamma$ , we were obliged to use a period between the detachment and the re-absorption of the eddy.

According to NP, the coefficient  $\alpha$  decreases quickly when its initial value is unity and the outflow PV is zero. In the case of finite PV outflow,  $\alpha$  decreases more slowly, starting from about 0.22. Our situation was analogous when we started with  $\alpha=1$  and  
25  $\alpha=0.4$ . However, in most of the experiments that started with  $\alpha=0.1$ , we observed a slight growth in  $\alpha$ , especially for large  $\gamma$ . Most likely, this was a consequence of a decrease in the size of the meander/eddy due to  $\beta$ . We only once obtained a very

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slight decrease in  $\alpha$  due to the combined effect of viscosity and a decreasing size. Resulting from the above-mentioned behaviors of  $\alpha$ , its average value accumulated in relatively narrow intervals: between 0.12 and 0.36 for  $h_0=0$ , and between 0.08 and 0.33 for  $h_0=300$  m. The scattering weakened as  $\gamma$  was increased.

Figure 8 shows a comparison of  $R_f$  between the analytics and the numerical simulations. Bearing in mind that the starting values of  $\alpha$  in the numerics were 1, 0.4, and 0.1, we mark the numerical results for the time-averaged values of  $\alpha$  with circles and diamonds. Analyzing Fig. 8 a–c, where  $\gamma$  is  $0^\circ$ ,  $30^\circ$  and  $60^\circ$  (respectively), we see that the numerical simulations confirm the theoretical tendency for the radius to decrease with growing angle. However, such a decrease is noticeably weaker than in the theoretical prediction. In this connection, the scattering in  $R_f$  also weakens with  $\gamma$ . For example, in the case  $h_0=0$  and starting from  $\alpha=0.1$ , 0.4, and 1.0, we obtained the values 381, 236, and 255 km for  $\gamma=0^\circ$  (represented as diamonds in the upper panel) and 203, 181, and 213 km for  $\gamma=75^\circ$  (not shown). The theoretically predicted ranges, corresponding to the above-mentioned intervals, are between 240 and 350 km for  $\gamma=0^\circ$ , and between 160 and 165 km for  $\gamma=75^\circ$ . Strictly speaking, for  $\gamma=75^\circ$ , we give here values corresponding to a greater  $\alpha$ . In our theory, we cannot calculate  $R_f$  corresponding to the  $\alpha$  obtained in the numerics when  $\gamma$  is large, because under such conditions, the BE is forced into the wall. The same situation caused the appearance of two numerical circles to the left of the theoretical curves in the lower plot ( $\gamma=60^\circ$ ). We conclude that, for nonzero  $\gamma$ , the numerical radii calculations give larger values than our theoretical model. However, the differences between the two are small. In the case of a zonal wall, the agreement is better overall, but the scattering of numerical values is significant.

Figure 9 shows a comparison of the theoretical and numerical values of  $C_\xi^z$  for  $\gamma=0^\circ$ ,  $30^\circ$ , and  $60^\circ$  whereas Fig. 10 shows the comparison of  $\Phi$  for the same angle. We note that there are two circles instead of three in the lower panel of Fig. 9. This is because, as mentioned, in our simulation with  $\gamma=60^\circ$ ,  $\alpha=1$  and  $h_0=300$  m we did not achieve detachment, so we could not compute the eddy drift velocity. On the whole, the agreement is satisfactory. Concerning  $C_\xi^z$ , we can say that this agreement obviously

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improves with growing angle  $\gamma$  for  $h_0=0$ , and in the case of  $h_0=300$  m, it looks slightly worse because of the noticeable scattering in the numerical values. Most importantly, the agreement of the theoretical and numerical eddy drift velocity confirms (3) which was introduced without any rigorous analysis.

5 The agreement seems to be worse with regard to  $\Phi$ , especially for nonzero  $\gamma$ . However, it is not easy to compute the mass flux numerically because of noise (e.g., in the forms of Kelvin waves and secondary meandering of the outgoing flow, and frictional effects) leading to ambiguity in the boundaries of this flux. The scattering of the numerical values ranges between slightly negative and about 0.18, which is admissible when  
10 the theoretical values are between zero and 0.15 (for  $\gamma=30^\circ$ ), or even between zero and 0.04 (for  $\gamma=60^\circ$ ). In either case, the numerics clearly confirm that  $\Phi$  is considerably smaller when the coastline is highly slanted.

## 5 Summary and conclusions

15 The main aim of our theoretical and numerical analyses was to examine the hypothesis that the glacial AC was similar to the present day EAC where, due to the orientation of the coastline, no mean rings shedding is usually observed. According to our hypothesis, the shedding of eddies is severely restricted when the slant angle  $\gamma$  is greater than  $65^\circ$  (the present EAC, and the glacial AC) but occurs steadily when  $\gamma$  is smaller than about  $20^\circ$  (present AC).

20 To examine this hypothesis we developed a non-linear model of retroreflecting currents that flow along slanted coastlines. We studied the dependence of rings diameter, speed and frequency of shedding, on the coastline slant, the PV of the formed eddies, and the thickness of the surrounding upper layer. The results are shown in Figs. 3–5 illustrating that there is significant shedding when the slant is small (sub-critical angle) but almost  
25 no shedding when the slant is large (supercritical angle).

Although we do not give an exact definition of the critical angle, we treated the angle of visible convergence of the lower and upper boundaries of the theoretical eddy radius

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(Fig. 3) as the critical angle. Indeed, according to our definition of the lower boundary (osculating rings), the convergence of the lower and upper boundaries means that a chain of detached rings must form downstream. This implies that, in practice, the rings are likely to hinder each other owing, for instance, to viscosity. Because for supercritical slants, small long-shore drift speeds are predicted for detached rings, the slowly moving eddy will be hindered and recaptured by the one behind it or by a meander of the retroflected current. Such a scenario agrees well with the dynamics retroflected current EAC rings described by Nilsson and Cresswell (1981), and Sokolov and Rintoul (2000). These rings usually stay at the same place for a long time and may eventually re-coalesce with the EAC.

We used a modified version of the Bleck and Boudra (1986) reduced gravity isopycnic model and obtained plots surprisingly similar to the observed EAC dynamics when the slant was taken to be  $60^\circ$  or more (Fig. 7). Also, as expected, we obtained chains of detached eddies when the slant was  $15^\circ$  (Fig. 6) and, a less clear chain when  $\gamma=30^\circ$ . The transition range of slant angles was between  $35^\circ$  and  $55^\circ$ , which agrees with our description of the critical angle that is approximately  $20^\circ$ ,  $40^\circ$ ,  $60^\circ$ , and  $65^\circ$  when  $\alpha$  is 0.1, 0.2, 0.5, and unity (respectively). During most of our simulation,  $\alpha$  was between 0.15 and 0.3 and the variables obtained in the theoretical modeling were in quantitative agreement with our numerical simulations (Figs. 8–10). On this basis, we suggest that the significant reduction in the exchange between the Indian and South Atlantic Oceans during the glacial and the YD was due to a northward migration of the WSC. This led the AC retroflection area to shift to a latitude of a supercritical coastal slant (Fig. 1). Other important results of our study can be summarized as follows:

- An increase in  $\gamma$  leads to a decrease in the radii of detached rings and makes them less sensitive to variations in  $\alpha$  (Fig. 3). Nevertheless, even in the case of supercritical slant, the theoretical values of  $R_f$  are 200–220 km and, therefore, still noticeably greater than the observational values.
- The mass-flux ratio  $\Phi$  decreases monotonically with increasing coastal slant

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(Fig. 5). When  $\gamma \geq 6^\circ$ , it becomes less than unity (even for eddies with zero PV, i.e., intense eddies) implying that the “vorticity paradox” discussed by ZN is circumvented.

- Our assumption of a nearly circular BE fails when the PV is large (weak eddies) and  $\gamma$  is not very small. In that case, the BE does not grow; rather, it is forced into the wall and deforms.

Concerning the distance between two consecutive eddies, we note that the ratio  $R_{fu}/R_{fl}$  is maximal in the case of zonal wall, when it is approximately 1.11 (i.e., less than  $2^{1/5}$ ). Therefore, we conclude that, theoretically, the separation distance should not exceed the eddy diameter. This conclusion is not far from the observational data, even though the eddy diameter is smaller than our theory yields. Indeed, if the eddies are shed on average 5 times per year, and their migration rate varies between 2 and 10 cm/s, then, during the period of generation, the eddy migrates 380 km on average, and at the most 630 km. Taking into account a typical eddy radius of 140 km, we obtain that the ratio  $(d/2R)$  is 0.36 in average, with a maximal value of 1.25.

As is frequently the case, our ability to compare our theoretical results with the numerical simulations was limited due to the effect of viscosity in the numerics, which led to a relatively narrow range of  $\alpha$ . In addition, the viscosity in the numerics makes the outgoing flux appear blurry, resulting in a possibility of errors as large as 0.2 in the determination of  $\Phi$ . Despite both of these aspects, our comparison is very useful. Although in our numerical simulations we confirmed the observation that capturing and re-splitting of eddies can be a possible cause of non-regularity in their shedding (Veldhoven, 2005), we find that non-regularity could also be connected with variability of the retroreflection position (Lutjeharms, 2006). For example, we note that Esper et al. (2004) pointed out seasonality of its position.

Our results also agree with Chassignet and Boudra’s (1988) sensitivity analysis, which showed that decreasing coastal slant leads to an increase in the production of rings. On the other hand, the numerical experiments by Pichevin et al. (1999) showed

that the dependence of the periodicity of rings shedding on the slant angle could be negligible. This might be a result of the specific geometry of coastline in their model. We leave this issue as a subject for future investigations. Finally, we note that taking into account the coastline slant could also improve understanding of other retroreflecting oceanic currents such as the North Brazil Current (NBC). Unfortunately, the question of a priori determination of the eddy PV remains unanswered, and this is significant particularly when the vorticity of the incoming fluid is cyclonic rather than anticyclonic.

## Appendix A List of symbols

AC	–	Agulhas Current
BE	–	basic eddy
$C_x$	–	eddy velocity in the open ocean
$C_\xi$	–	eddy migration rate along the slanted coast
$C_{\xi f}$	–	eddy migration rate after detachment
$C_{\xi l}, C_{\xi u}$	–	values of $C_{\xi f}$ for eddies with radii $R_{fl}, R_{fu}$ , respectively
$d$	–	distance between consecutive eddies
$d_u$	–	“upper” boundary of $d$
EAC	–	East Australian Current
$g'$	–	reduced gravity
$h_0$	–	upper layer thickness at the wall
$H$	–	upper layer thickness outside the retroreflection area
$H_{fl}, H_{fu}$	–	values of $H$ at the moments $t_{fl}, t_{fu}$ , respectively
LGM	–	Last Glacial Maximum
MOC	–	meridional overturning circulation
NADW	–	North Atlantic Deep Water
NP	–	Nof and Pichevin (2001)
PV	–	potential vorticity

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$Q$	– mass flux of the incoming current
$q$	– mass flux of the retroflected current
$R$	– radius of the eddy (a function of time)
$R_i$	– initial radius of the eddy
$R_f$	– radius of detached eddy
$R_{fl}, R_{fu}$	– “lower” and “upper” boundaries of $R_f$
$t$	– time
$t_f$	– period of the eddies generation
$t_{fl}, t_{fu}$	– “lower” and “upper” boundaries of $t_f$
WSC	– wind stress curl
YD	– Younger Dryas
ZN	– Zharkov and Nof, paper submitted to “Ocean Sciences”
$\alpha$	– vorticity (twice the Rossby number)
$\beta$	– meridional gradient of the Coriolis parameter
$\gamma$	– slant of coastline
$\xi, \eta$	– axes of rotated moving coordinate system
$\Phi$	– ratio of mass flux going into the eddies and incoming mass flux
$\Phi_l, \Phi_u$	– values of $\Phi$ for eddies with radii $R_{fl}, R_{fu}$ , respectively

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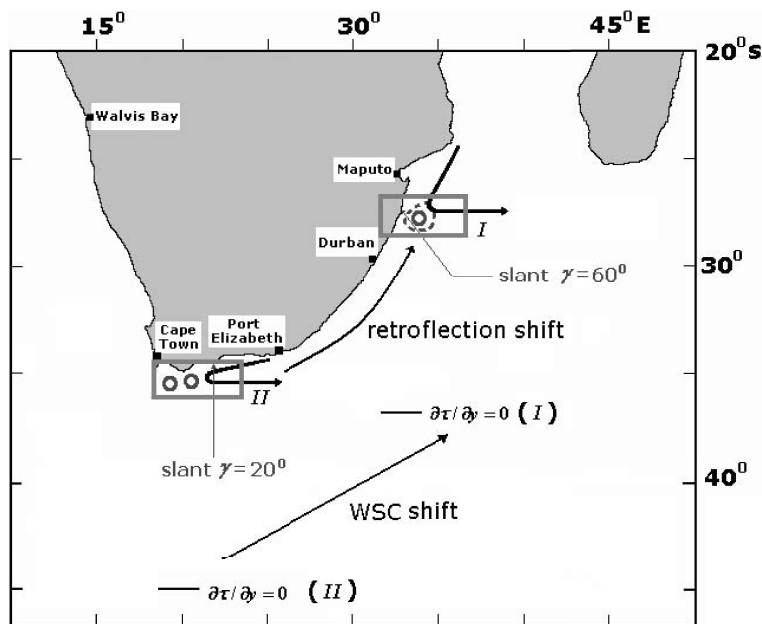
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**Fig. 1.** Schematic diagram of the Agulhas retroflexion and the detached rings. Note that the slant of the coastline relative to the meridional direction varies dramatically as one moves northward along the coast. Box I displays glacial conditions: wind stress curl vanished at the lower latitudes where the coastal slant is about  $60^\circ$ . Ring shedding was rare because the translation velocity of the detached rings along the wall was small. Some rings could be dissipated or reabsorbed by the meandering current. Box II displays post-glacial conditions: wind stress curl vanishes at higher latitudes where the coastal slant is about  $20^\circ$ . A chain of rings is regularly shed because the migration speed along the wall is high.

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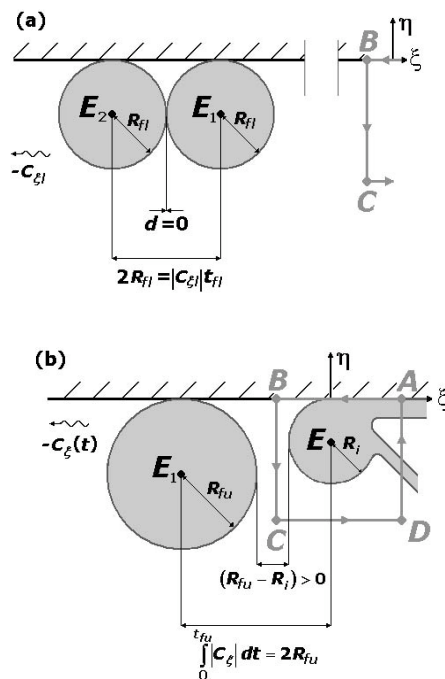
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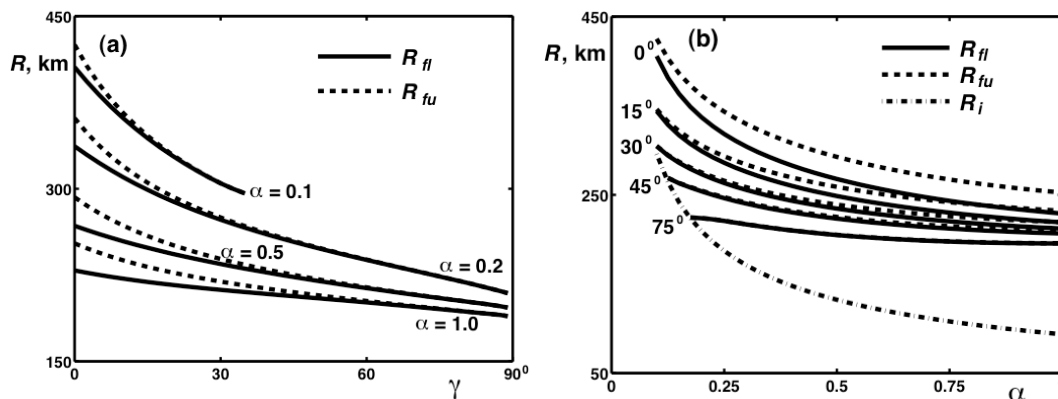


**Fig. 2.** Geometries associated with the lower and upper boundaries of the final eddy radius. The upper panel **(a)** shows two consecutive osculating eddies ( $d=0$ ) away from the retroflection. The detachment period  $t_{fi}$  is obtained by a division of the doubled-finaleddy- radius,  $2R_{fi}$ , by the modulus of the final eddy migration rate. The segment BC corresponds to the western boundary of the integration area (see ZN). The lower panel **(b)** shows the already detached eddy (centered in  $E_1$ ) migrating at the rate  $-C_s(t)$ , and an incipient basic eddy (BE) centered in E. At the moment  $t_{fu}$ , the distance between the two eddies is  $(R_{fu} - R_i)$ , which is positive because the incipient BE is less developed. Therefore, our ABCD contour encloses only the incipient BE.

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**Fig. 3.** The theoretical solutions. **(a)**  $R_{fl}$  (solid lines) and  $R_{fu}$  (dashed lines) plotted against the angle  $\gamma$ . Each pair of convergent curves is marked by a corresponding value of  $\alpha$ . **(b)**  $R_{fl}$  (solid lines),  $R_{fu}$  (dashed lines), and  $R_i$  (dash-and-dotted line) plotted against  $\alpha$ . Each pair of divergent curves is marked by a corresponding value of  $\gamma$ . It is seen that the curves corresponding to  $\gamma \geq 45^\circ$  start from the curve of  $R_i$ . Note that lines depicting  $R_{fl}$  and  $R_{fu}$  for  $\gamma = 75^\circ$  almost coincide. Also, note that, as should be the case, the upper boundary is always above the lower boundary. In both plots:  $h_0 = 0$ ,  $\beta = 2.3 \times 10^{-11} \text{ m}^{-1} \text{ s}^{-1}$ .

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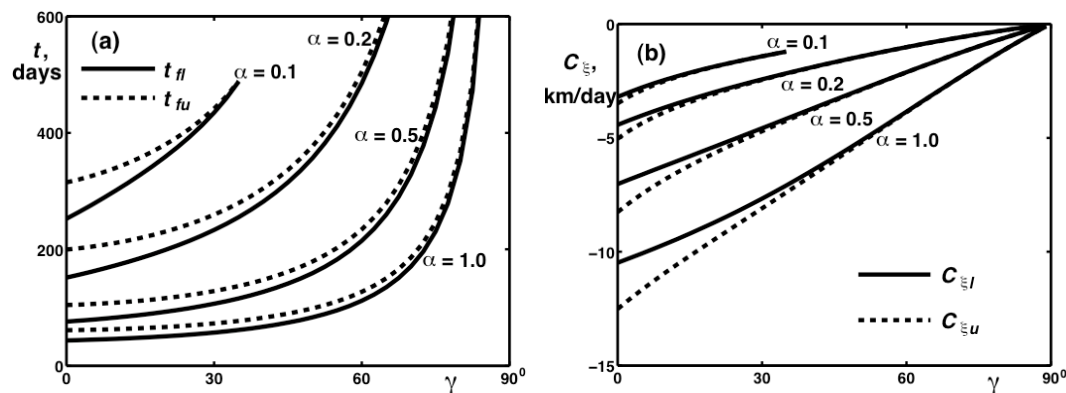
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**Fig. 4.** As in Fig. 3, but for the detachment period and migration rate. **(a)** Graphs of  $t_{fl}$  (solid lines) and  $t_{fu}$  (dashed lines). **(b)** Graphs of  $C_{\xi l}$  (solid lines) and  $C_{\xi u}$  (dashed lines) against the angle  $\gamma$  for the same values of  $\alpha$  as in Fig. 3a. The dashed line is above the solid line (left panel) and below the solid line (right panel) as should be the case.

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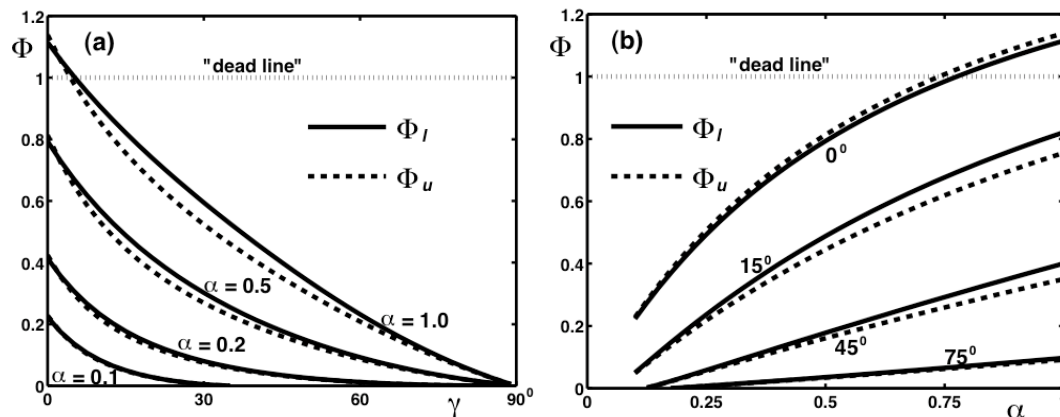
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**Fig. 5.** Theoretically based plots of  $\Phi_I$  (solid lines) and  $\Phi_u$  (dashed lines): **(a)** against the angle  $\gamma$  for the same values of  $\alpha$  as in Fig. 3a; **(b)** against  $\alpha$  for the same values of  $\gamma$  as in Fig. 3b. The “dead lines” are the same as in Fig. 5 of ZN. As before,  $h_0=0$ , and  $\beta=2.3 \times 10^{-11} \text{m}^{-1} \text{s}^{-1}$ .

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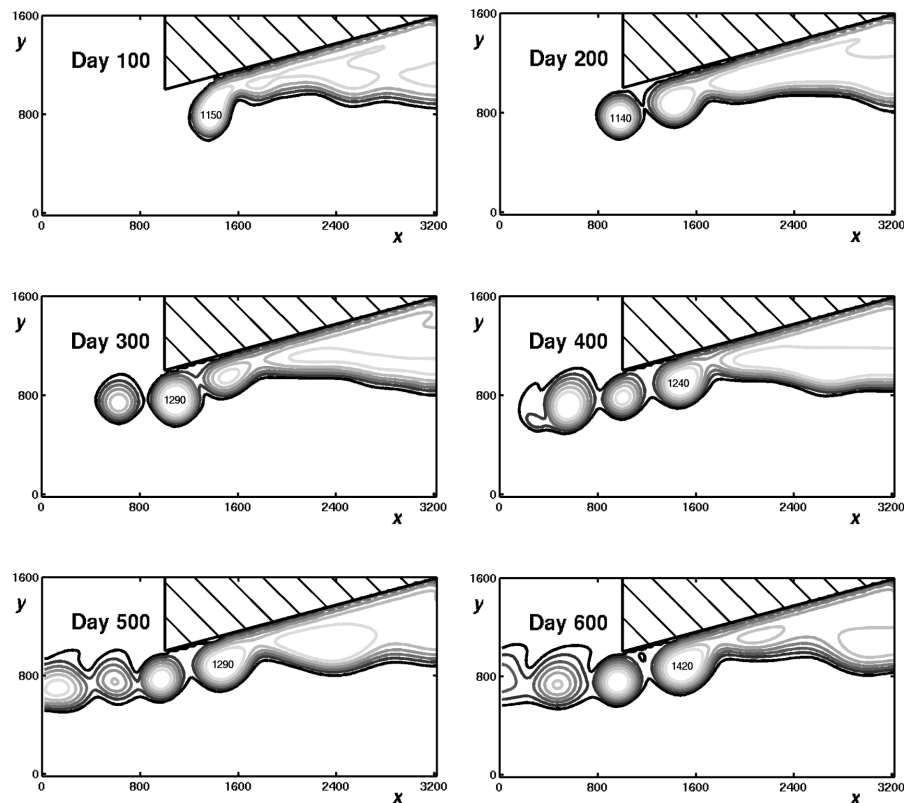
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**Fig. 6.** Thickness contours of the numerical simulation for the first 600 d. A chain of eddies is formed despite intermittent cases of eddies recapturing. Numerical values:  $\gamma=15^\circ$ ,  $\alpha=1$ ,  $h_0=0$ , and  $\nu=1000 \text{ m}^2 \text{ s}^{-1}$ . Spacing between contours represents increments of 200 m, and the maximal thickness of the upper layer is given in meters. The x and y scales are in kilometers. Note that we used  $\beta=6 \times 10^{-11} \text{ m}^{-1} \text{ s}^{-1}$  here and in all the following figures.

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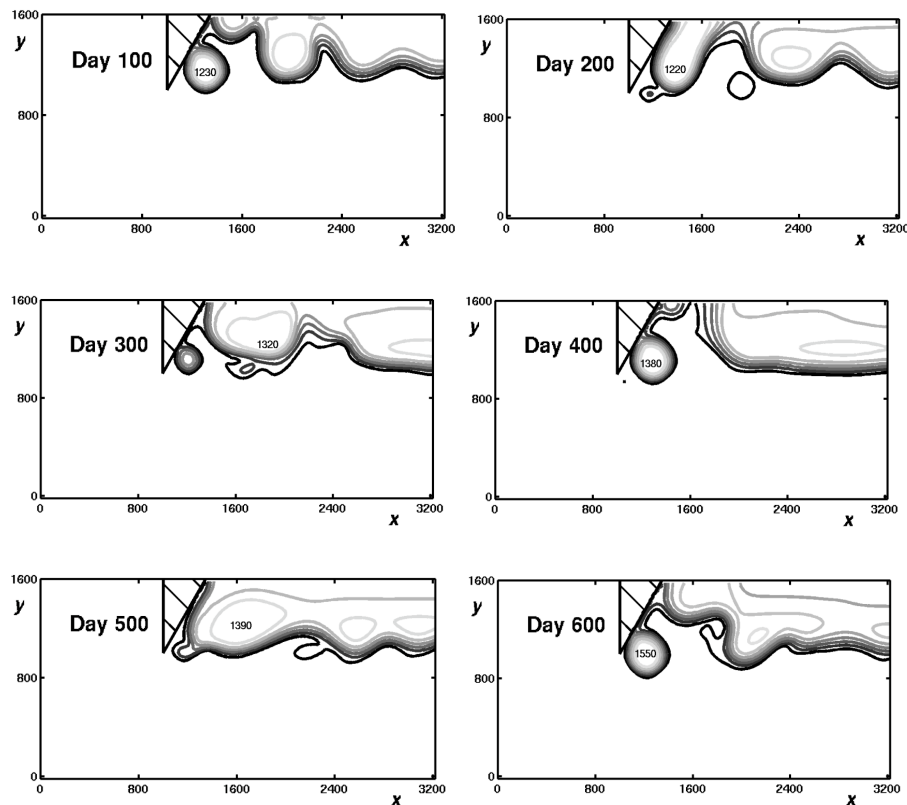
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**Fig. 7.** The same as Fig. 6, but for  $\gamma=60^\circ$ ,  $\alpha=1$ ,  $h_0=0$ , and  $\nu=700 \text{ m}^2 \text{ s}^{-1}$ . The initially detached eddy is re-captured by the incipient meander. The merged eddies then detach but are recaptured again later; this process repeated itself 3–4 times. The eddy shown in day 600 was later re-captured by day 700 but this recapturing is not displayed in the figure.

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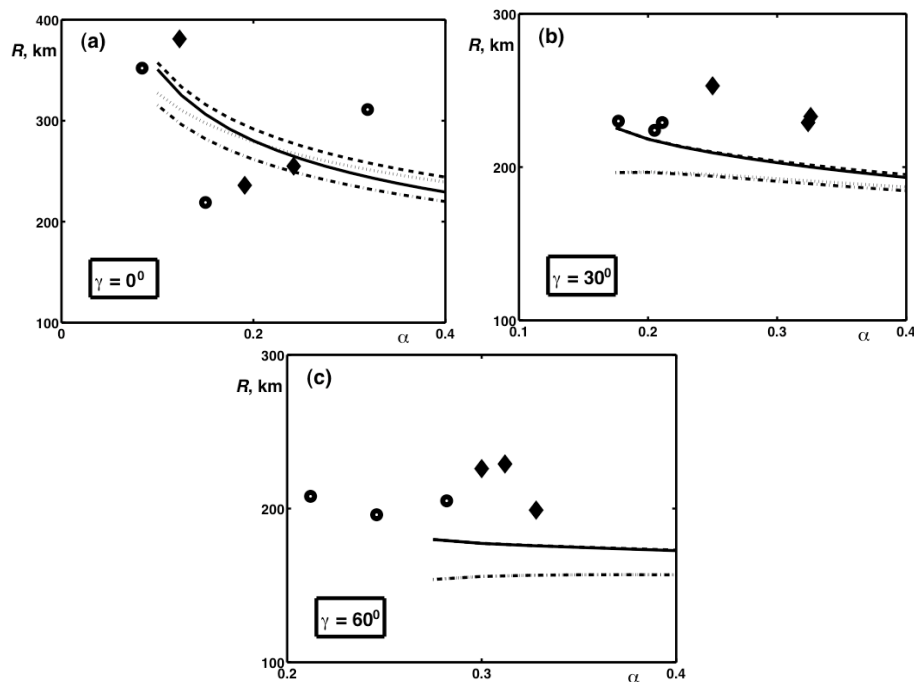
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**Fig. 8.** Comparison of modeled radii with the numerics. The modeled values of  $R_{fI}$  and  $R_{fU}$  are plotted against  $\alpha$  for  $h_0=0$  (solid and dashed lines, respectively) and for  $h_0=300$  m (dash-and-dotted and dotted lines, respectively). The numerical values of  $R$  correspond to  $\alpha$  averaged over the time of the experiments (diamonds for  $h_0=0$  and circles for  $h_0=300$  m. Here: **(a)**  $\gamma=0^\circ$ , and the theoretical curves start from  $\alpha=0.1$ ; **(b)**  $\gamma=30^\circ$ ; **(c)**  $\gamma=60^\circ$ . The theoretical curves start from the points at which  $R_{fI}=R_i$  or  $R_{fU}=R_j$ . We note that one pair of joined circles and one pair of joined diamonds are also seen in Fig. 8b (even though the parameters have different values). The curves of  $R_{fI}$  and  $R_{fU}$  are overlapping in the lower panel because the angle of  $60^\circ$  is supercritical.

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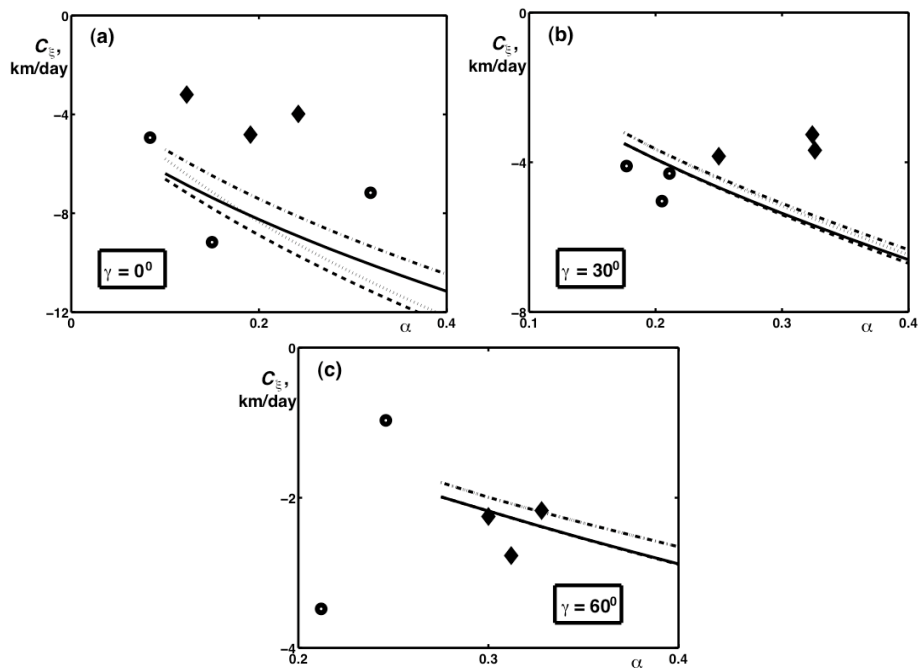
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**Fig. 9.** As in Fig. 8 except that this is a comparison of the modeled propagation rates with the numerics.

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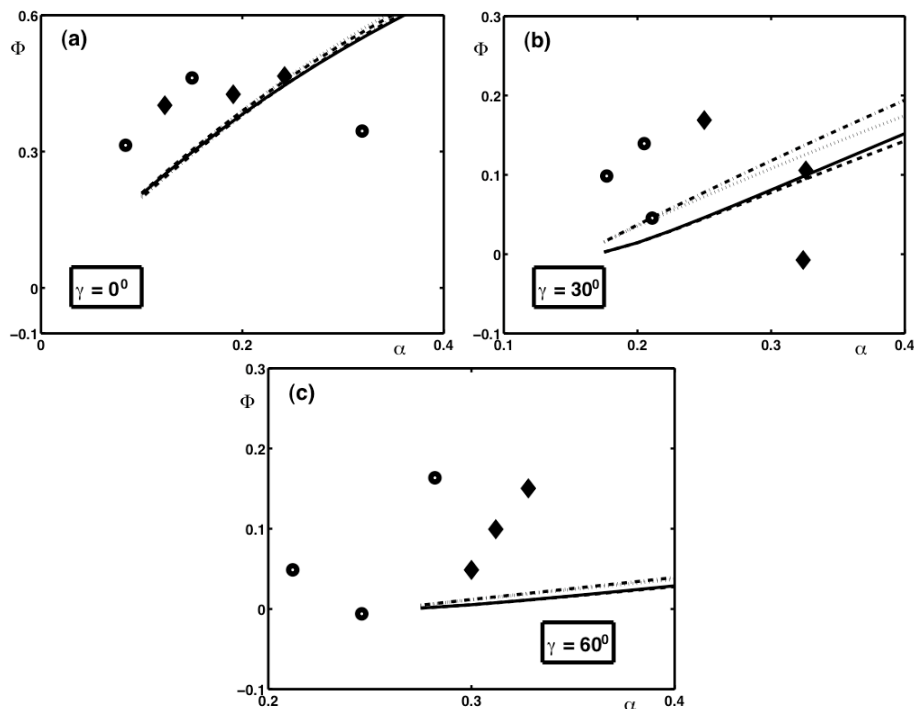
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**Fig. 10.** As in Fig. 8 except that we show here a comparison of the mass flux ratio  $\Phi$ .

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