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Thermal resonance of the atmosphere to SST anomalies. Implications for the Antarctic circumpolar wave

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ABSTRACT

It is well-known that when Rossby waves are stationary in the belt of mid-latitude westerlies, resonance conditions occur allowing the atmospheric response to external perturbations to be greatly enhanced. The concept lies at the heart of the interaction of planetary flows with topographic mountain chains. In contrast, early studies of the atmospheric response to thermal forcing, focussed on the off resonance response. Now available, many GCM studies dealing with the atmospheric response to prescribed SST have revealed a great variety of responses, the difficulty of extracting the signal at planetary scale being compounded by the overwhelming activity of transient eddies at the synoptic scale. Given the long lasting influence of large scale SST anomalies, climate predictability may be expected to improve if one can identify feedbacks between the large scale climate anomalies and the SST distribution. We propose a simple theory that explores the physics of this atmospheric response to SST and may suggest ways to analyze data from the more complex GCM. We neglect all interactions between the transient eddies and the large scale waves that would go beyond Fickian mixing laws although there is evidence that the transient eddies by themselves take part in the maintenance of the low frequency variability. Because the observed perturbations are small, a linear theory is appropriate. Using a 2-level model of the atmosphere, a resonance condition occurs when the Rossby waves are stationary against the vertically averaged mean zonal flow. The resonance is sharp when eddy dissipation through surface friction is small. In a small wavenumber window controlled by the vertical structure of the mean flow, the response is equivalent barotropic and baroclinic elsewhere. Only in this window is the familiar response of high SLP downstream of warm SST recovered. For certain combination of thermal damping and surface drag, the atmospheric response is amplified to produce a positive feedback on the SST. When the atmospheric model is coupled to a one and a half level ocean model with a zonally periodic geometry appropriate to the Southern Oceans, a linear instability appears. The application of this process of thermal resonance to the Antarctic circumpolar wave is discussed. We find that under realistic values of the mean state, an unstable coupled wave emerges in a narrow wavenumber window which coincides with the near resonance conditions found previously in the atmospheric model. This instability provides a powerful scale selective mechanism for the perturbations. The thermal forcing is primarily responsible for amplification of the SST while the powerful Antarctic circumpolar current is primarily responsible for advecting the anomalies around the globe and setting the period.

1. Introduction

There is increasing evidence from the analysis of historical data sets that co-variation of SST (sea surface temperature) and SLP (sea level pressure) exist at sub-basin scales. Palmer and Zhaobo (1985) carried out a composite analysis of a 30-year period and found that warm SST anomalies off Newfoundland are associated with pressure highs lying downstream (eastward) of the SST anomaly. With longer records becoming available,

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the structures were shown by Kushnir (1994) and Kushnir and Held (1996) to be frequency dependent. At interannual periods, the atmosphere is apparently close to an equivalent barotropic structure with high (low) SLP anomaly associated with warm SST anomaly while at interdecadal periods, a dominant low is found east of the warm anomaly indicative of a baroclinic response. In the North Pacific the pattern of high SLP east of warm SST occurs at both interannual and decadal scales, Tanimoto (1993), Deser and Blackmon (1995). In the Southern Hemisphere although the records are shorter a remarkable phase relationship of warm (cold) SST associated with high (low) SLP (leading eastward with a nearly 90° phase advance) has been discovered by White and Peterson (1996). The Antarctic circumpolar wave (ACW) is periodic with a dominant wavenumber 2 around the globe propagating eastward at 56°S at speed of 6–8 cm s^{-1}, with a period of 4–5 years. Presumably the simpler channel geometry of the southern mid-latitudes (with zonal communication between ocean basins and absence of topographic obstacles in the atmosphere) is an important factor for the existence of this coherent planetary structure.

The question of the planetary wave response of the atmosphere to prescribed SST (and to heating in general) is a long-standing one which has been addressed through both theoretical studies and GCM's analysis. The former were initiated by Smagorinsky (1953) who sought to find out if thermal forcing could compete with topographic forcing (Charney and Eliassen, 1949) to explain the existence of stationary waves. Given the observed amplitudes of the stationary wave response, Smagorinsky explored the linear quasigeostrophic theory of perturbations of a mean zonal flow with constant vertical shear to anomalous heating. He found that away from resonance the thermal effect was usually smaller than the orographic effect. The analytical complexities of the continuous model prompted Derome and Wiin-Nielsen (1971), Egger (1977), Chen and Trenberth (1985) and many others to reexamine the problem with a 2-level model. Egger's comparisons with Smagorinsky's results were specially good off resonance. Numerical computations (with a GCM) in such regimes showed that a downstream low was associated with a positive mid latitude SST anomaly. Hoskins and Karoly (1981) used perceptive scaling analysis to describe the variety of atmospheric response to heating. Their description at planetary scale is the following: when meridional heat transport by anomalous equatorward winds balance a positive heating anomaly, a pressure low is found east of the heating. The vorticity balance is one between low level divergence (negative stretching) and planetary vorticity advection. In this situation the pressure anomaly changes sign at some intermediate level and the response is baroclinic. The comparison of 2-level, linear, quasigeostrophic model and 2-level GCMs' away from resonance, Roads (1980), was generally good with the baroclinic response of warm lows downstream of heating regions. These essentially forced responses to thermal forcing were reviewed by Held (1983) and Frankignoul (1985). The interesting possibility of a near resonance (so that free Rossby waves can be excited) has been discussed by Tung and Lindzen (1979) in the context of the initiation of stratospheric warmings from lower tropospheric levels. Roads (1982) pursued Smagoringsky's study using WKB methods and stressed that resonance could alter significantly previous conclusions. Surface friction was introduced and shown in conjunction with thermal damping to amplify the response near resonance.

Spurred by the objective of improving seasonal predictability, the response of many GCMs to prescribed SST have been obtained. A recent review by Peng et al. (1997) reveals the diversity of the results. Low resolution model simulations range from no response to baroclinic to equivalent barotropic responses. High resolution experiments such as that of Palmer and Zhaobo (1985), exhibited equivalent barotropic structure but with amplitudes somewhat on the weak side as compared to the observations. They proposed a theoretical argument based on a 2-level model and showed that the response to a mid-level tropospheric heating could be equivalent barotropic at a wavenumber making the disturbances stationary with respect to some average of the upper and lower layer mean zonal flow. The low level convergence induced a poleward flow over the source via the Sverdrup balance while advection of relative vorticity important in the upper layer helped to keep the flow in the same direction at upper levels, so that high pressures were found downstream of the heating at each level. They discussed furthermore how this wind distribution could create a positive feedback on the SST through heat transport by anomalous Ekman velocities in the oceanic mixed layer. As will become apparent, the present work is a direct outgrowth of their work. Carrying out similar experiments, Peng et al. (1995) recovered Palmer and Zhaobo's response under November conditions but found just the opposite for January conditions and indicated that "the different sensitivities to the midlatitude SST anomalies may have resulted, in part, from the differences in their background circulation states". This issue was pursued by Peng et al. (1997) and Peng and Whitaker (1999). The latter work showed that the vorticity fluxes by the transient eddies were important in forcing the anomalous pressure ridge found downstream of the positive SST in the perpetual February conditions but not in the perpetual January conditions.

Given this somewhat diverse picture of the atmospheric response to SST and our objective to isolate important ocean-atmosphere feedbacks at planetary scales, we did not feel confident to follow the approach of simply forcing the equivalent barotropic of the atmosphere by requiring surface pressure to be proportional to surface temperature a choice made by Qiu and Jin (1997) in the context of the dynamics of the Antarctic Circumpolar wave, Liu (1993) and Talley (1999). Instead Goodmann and Marshall (1999) chose for this problem a quasi geostrophic 2-layer model of the atmosphere similar to what is used herein.

Both Qiu and Jin (1997) and Goodman and Marshall (1999) coupled their model atmosphere to ocean models composed of a surface layer and one or two lower layer and proposed an unstable coupling mechanism for the Antarctic circumpolar wave through which the anomalous winds produced say by the pressure high downstream of a positive anomaly induce a downward Ekman Pumping in the ocean that re-inforce the initial temperature anomaly. White et al. (1998) proposed a similar feedback but based on meridional anomalous heat transport by Ekman velocities. All these previous works rely on anomalous "mechanically induced" meridional oceanic heat fluxes to release the potential energy stored in the oceanic thermocline. The proximity of a Rossby wave resonance immediately comes to mind as a possible scale selective mechanism and indeed this choice is made at the outset in the Goodman and Marshall equilibration mechanism. Given the fast time scale response of the atmosphere, a prediction of that scale is tantamount to a prediction of the period since the anomalies are transported at the slow rate of oceanic speed with the atmosphere instantaneously adjusting to the new positions of the anomalies. White and Peterson's observations suggest strongly that the scale of the response is selected dynamically and not by external scales linked to ocean-continent distribution that would favor stationary over travelling patterns. (The situation is far more complex in the Northern Hemisphere with the North Atlantic Oscillation because additional anomalous forcing by the land–sea temperature contrast competes with the forcing by travelling SST anomalies.)

We shall propose a physical analog whose objectives are two fold, first explore the linear response of the atmosphere to SST forcing and second describe a new process through which an unstable coupled pattern (similar to what is known about the Antarctic circumpolar wave) may emerge dynamically. The physical analog must be simple enough to be tractable analytically and complete enough to be relevant. The linear theory of the atmospheric response to a periodic SST distribution is reexamined following the Palmer and Zhaobo suggestions to keep advection of relative vorticity. Talley (1999) estimated that $U/\beta L^2$ (~ 0.02) was small for the ACW using a wind velocity U of 10 m s^{-1} and a length scale L of 6000 km (half the wavelength of the ACW). This underestimates relative vorticity because anomaly patterns (specially those from the Southern hemisphere) have meridional scales at least twice shorter than zonal scales so that the above estimate now multiplied by 5 becomes 0.1, small but not very small. The atmospheric model described in Section 2 has two vertical levels. Phillips (1951) proposed that such a low resolution might provide adequate answers if the vertical velocity field has a single extremum across the troposphere. Although he had in mind the synoptic development of weather systems, Egger (1977) used this idealization at planetary scales with favorable comparisons with higher resolution models. This type of discretization falls short of taking into account the observed facts that stationary wave activity penetrates beyond the tropopause well into the stratosphere. Putting a rigid lid at the tropopause may therefore seem a serious oversimplification. Furthermore the lid can have serious adverse effects by adding false resonances Tung and Lindzen (1979). Given the value of Phillips's model in the baroclinic instability problem and the impossibility to solve analytically the vertical structure equation when a realistic basic state is introduced, several works have analysed the effect of the rigid lid approximation on the planetary waves in the 2-level model and given the conditions under which the spurious resonance can be avoided. Chen and Trenberth (1985) first showed that the upper boundary conditions (radiating or rigid lid) have minor impacts when the vertical shear and lower level winds are large enough. There is of course numerical differences in the precise value of the external Rossby mode resonance with each type of boundary conditions but the response to thermal (and orographic forcing) are quite similar. Panetta et al. (1987) compared the two layer (equal depth) modes against the continuous Charney and hyperbolic tangent profiles modes and found a close physical correspondence when the shear is supercritical in the sense of Phillips's baroclinic instability criterion. With lower layer westerly winds they showed that this implied that the vertically averaged zonal mean flow had to be greater than $\beta\lambda^2$ where λ is the internal Rossby radius of deformation. Only the external Rossby mode may then propagate horizontally and no resonance appears with the internal mode. Our study is focussed on the latitudes of the strong mid-latitudes westerly jets and carried out in a parameter regime where no such internal spurious resonance happens. A further simplification to keep the problem tractable analytically is to neglect wave dispersion caused by the meridional structure of the mean flow. This point is reviewed by Held (1983). If a near resonance condition is invoked as a scale selective mechanism the existence of turning points in latitudes helps to build the wave guide. However the critical layer (the zero wind line in the tropics in the stationary wave case) is of a nature "absorbing or reflecting" that depend on the local dynamics at the critical layer. Furthermore the 2-level model is a tool which is less appropriate in such a study for the reasons just given: spurious resonance will contaminate the solution in regions of weaker winds equatorwards of the westerly jets. The theory that we propose is therefore implicitly dependent on a weak leakage of wave energy from mid-latitudes to the Tropics through this critical layer.

Special attention has been given to guarantee energy conservation of the physical analog that we propose. This is specially desirable for comparisons with GCMs results. We consider mechanisms that are associated with very small growth rates (of the order of cycles per year) so that any energetic inconsistencies would cause much uncertainties in the results. It is also found that the vorticity dissipation by surface drag is absolutely crucial to the theory departing in this respect from the work of Goodman and Marshall (1999). This is because in the vicinity of resonance, higher order effects such as these become the dominant terms in the equations and the sensitivity of the large scale response to these higher order effects become important. Concerned with the interactions between the zonal flows and the large scale waves in each fluid, we do not consider explicitly the rôle of the transient eddies that stir the fluid vigorously at synoptic scales. It can be easily accommodated in the present theory whenever their effect is akin to Fickian diffusion of the large scale thermal anomalies. Indeed we can check a posteriori that the unstable growth rates of our coupled instability are large enough to survive to observed levels of eddy diffusivities. On the other hand, it is far more difficult (although very important) to devise a way of parameterizing such interactions at the vorticity level. The only vorticity interaction with unresolved turbulent scales that we have included is that due to surface drag. In more realistic models the transient eddies through their momentum fluxes are the primary agents for maintaining the low frequency anomalies in the GCM analyses of Branstator (1992). The interaction works both ways as Branstator (1995) shows that the storm track activity is itself structured by the circulation anomalies caused by the low frequency waves. The GCM response to Pacific SST from Peng and Whitaker (1999) points out similarly to the overwhelming importance of eddy feedbacks in the diversity of the large scale response to SST anomalies, response which is sensitive to the relative position of the surface heating and the storm track. The present ocean-atmosphere coupled mid-latitude instability which lacks these non Fickian dynamical feedbacks must therefore be looked at as no more than an elementary step in a hierarchy of studies aimed at discovering the long term climate predictability linked to the persistence of SST anomalies. Section 2 describes the atmospheric model and discusses the response to SST forcing. Section 3 presents the oceanic model and its coupling to the 2-level atmospheric model. It consists of a single active surface layer of an ocean moving at a prescribed mean velocity. The coupled model allows to compare the various feedbacks involving the wind stress effects on heat transport against direct thermal forcing that operates near resonance. Applications to the mid-latitude Southern Oceans are given in Section 4. We shall show that the energy source for the instability is the available potential energy associated with the westerly jet and that the main feedback between ocean and atmosphere is the surface heat flux, two results that depart from previous works.

2. The atmospheric response to SST

The objective here is to find out the conditions of existence of positive feedbacks on SST perturbations of a prescribed zonal mean flow. A linearized theory is justified, at planetary scales because waves are observed to be modest perturbations of the mean flow (in contrast to the prevailing situation at the synoptic scale). Restricting the interest to mid-latitudes cases, the quasi-geostrophic approximation is used and the mean zonal flow varies only in the vertical direction. Compressibility effects are handled through the use of log (pressure) coordinates. In this part, the focus is on the steady state response of the atmospheric to persisting SST anomalies in the form of Fourier components of given zonal wavenumber. To keep as close as possible to the discretization used in GCMs, Lorenz's 2-levels grid is used in the vertical, with horizontal velocities, pressure and temperatures carried out in the middle of levels and vertical velocity at the level boundaries (Fig. 1). Such a truncation and rigid lid upper boundary condition are thought to give correct physical responses provided that the mean zonal wind is strong enough to prevent spurious internal resonance (see the Introduction). The levels heights are defined through the use of σ coordinates, $p = p_0\,e^{-\sigma/H_0}$, with p_0 a constant surface pressure and H_0 the atmospheric scale height. With these σ coordinates the hydrostatic equation read:

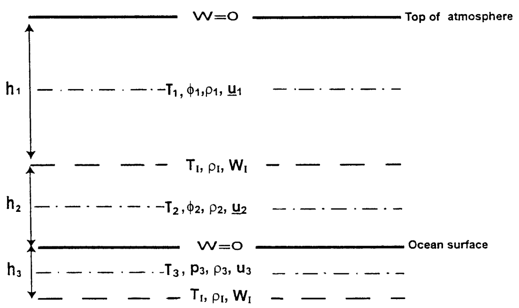


Fig. 1. The vertical discretization of the ocean- atmosphere model. The atmosphere is composed of two levels in σ − [log(pressure)] coordinates. The ocean is composed of a single active level for the perturbations. It is formulated in z coordinates.

$$\frac{\partial \phi}{\partial \sigma} = \frac{R}{H_0}\,T, \qquad (1)$$

where ϕ denotes the geopotential and R the perfect gas constant. Density variations in the continuity equation are restricted to vertical dependence, so that:

$$\rho \nabla_H \cdot \boldsymbol{u} + \frac{\partial}{\partial \sigma}(\rho w) = 0, \qquad (2)$$

where w is simply $\dot{\sigma}$ and, $\rho = \rho_a\,e^{-\sigma/H_0}$, ρ_a being the density at the ground. For analytic simplicity, a β plane approximation will be made neglecting metric terms in the continuity and momentum equations. With this approximation, the zonal x-coordinate stands for $R_T \cos\theta_0$ (longitude), with θ_0 a fixed reference latitude, R_T the radius of the earth, and the meridional y coordinate is $R_T(\theta - \theta_0)$. Because all departures from axisymmetry are neglected in the basic state, the perturbations variables are only advected by the zonal mean flow and any y-dependence is parametrical. One must keep in mind that there is an integer number of perturbations wavelengths around a latitude circle so that the only allowed wavenumbers k are $m/R_T \cos\theta_0$ with m an integer.

The horizontal momentum perturbation equations of a prescribed zonal mean flow U with vertical shear only are:

$$U\frac{\partial \boldsymbol{u}}{\partial x} + f\boldsymbol{k} \times \boldsymbol{u} = -\nabla\phi - \varepsilon \boldsymbol{u}, \qquad (3)$$

$\boldsymbol{k}$ being a vertical upward unit vector. The last term in (3) represents linearized turbulent friction at the ground. The last equation is the thermodynamic energy equation that reads in σ coordinates as:

$$U\frac{\partial t}{\partial x}+v\frac{\mathrm{d}\bar{T}}{\mathrm{d}y}+\left(\frac{\mathrm{d}\bar{T}}{\mathrm{d}\sigma}+\frac{K}{H_0}\bar{T}\right)w=\frac{1}{\rho C_\mathrm{p}}\frac{\partial Q}{\partial\sigma}; \tag{4}$$

$\mathrm{d}\bar{T}/\mathrm{d}y$ is the mean meridional temperature gradient associated geostrophically with the mean vertical shear, while $\mathrm{d}\bar{T}/\mathrm{d}\sigma+(K/H_0)\bar{T}$, the sum of vertical and adiabatic temperature gradient, measures static stability ($K = R/C_\mathrm{p}$). Because the steady state response is of interest here, the time derivatives have been neglected in (3) and (4). This has the effect of filtering out unstable baroclinic waves. Time dependence will be recovered on long oceanic time scales when coupling to the ocean model (Sections 3 and 4). We neglect latent heating (dry atmosphere) and restrict our attention to the effect of bottom heating of the lower layer in which Q is simply parameterized as a Newtonian flux $Q=\gamma(\mathrm{SST}-t)$ where SST is a prescribed periodic function, the objective being to calculate the atmospheric response to this perturbation. Finite differencing in the vertical implies a certain form of discretization of the temperature equation (4) and hydrostatic relation (1) if energy conservation properties and temperature variance conservation are enforced (Arakawa and Lamb, 1977). In this problem of slowly evolving large-scale waves, such requirements are very important to follow to avoid spurious growth rates. It is shown in Section 6, that if h_1 (h_2), ρ_1 (ρ_2) are, respectively, thickness and densities of upper (lower) levels, density and temperature are evaluated at the intermediate level as:

$$\rho_\mathrm{I}=\frac{(\rho_1 h_1+\rho_2 h_2)}{(h_1+h_2)}$$

and

$$t_\mathrm{I}=\frac{1}{2}(t_1+t_2).$$

With these choices, the discretized version of the perturbations eqs. (1–4) can be written as:

$$\phi_1-\phi_2=\frac{R}{H_0}t_\mathrm{I}\frac{(h_1+h_2)}{2}, \tag{5a}$$

$$U_i\frac{\partial\xi_i}{\partial x}+\beta v_i=(-1)^i\frac{f w_\mathrm{I}}{h_i^*}-\varepsilon\delta_{i2}\xi_i, \tag{5b}$$

$$U_i\frac{\partial t_i}{\partial x}+v_i\frac{\overline{\mathrm{d}T_i}}{\mathrm{d}y}+\left(\frac{\Delta T}{h_i^*}+\frac{T_\mathrm{I}K}{2H_0}\right)w_\mathrm{I}$$
$$=+\lambda\delta_{i2}(\mathrm{SST}-t_i), \tag{5c}$$

where eq. (5b) has been obtained after taking the curl of (3).

A rigid lid is present at the top of the atmosphere and the ground is flat. δ_{i2} is the Kronecker symbol and ε represents the turbulent drag at the ground. ΔT is $\frac{1}{2}\overline{(T_1-T_2)}$ and $\lambda=\gamma/\rho_\mathrm{a}C_\mathrm{p}h_2$. The symbols t_i and ξ_i indicate respectively perturbation temperature and relative vorticity. We see that in σ coordinates, equation (5b) looks like the incompressible form, once the new thicknesses h_i^* that appear are defined as:

$$h_i^*=\frac{h_i\rho_i}{\rho_\mathrm{I}}.$$

The quasi-geostrophic approximation can now be made, the streamfunction ψ_i being defined as ϕ_i/f_0 and perturbation velocities $\boldsymbol{u}_i$ as $\boldsymbol{k}\times\nabla\psi_i$. With this "closure", the set of equations (5) allows to evaluate perturbation velocities and temperatures given a Fourier component of SST. In practice, SST is written as the real part of ($\mathrm{SST}_0\,\mathrm{e}^{ik(x+qy)}$) and all variables are expanded similarly. The parameter q allows a possible dependence on the aspect ratio, meridional versus zonal, of the perturbations. The two vorticity equations are first used to eliminate the vertical velocities, providing a relation between ψ_1 and ψ_2. Along with (5a) this allows us to express ψ_1, ψ_2 and w_I as functions of t_1 and t_2. By inserting these relations into (5c), one obtains two linear equations containing the two temperature variables, which can be solved in terms of the amplitude SST_0.

Because the unforced set of equations allows stationary Rossby wave in case of westerly flows, resonance conditions are expected for the external Rossby mode. When the system is close to resonance, there may be situations in which the temperature response of the lower layer is not too out of phase with the SST and larger than the SST. If we find situations in which this happens, we shall have demonstrated the existence of a positive feedback on the SST. We explore this possibility and judge the efficiency of this feedback later via a coupled atmosphere–ocean model later.

The external mode resonance conditions can be found by expressing the streamfunctions in terms of the temperatures. Let us define the baroclinic mode for the streamfunction amplitudes $\psi^- = \psi_1 - \psi_2$ and eliminate w_I between the 2 vorticity equations (5b):

$$\psi^- = \frac{R}{f_0 H_0}\left(\frac{h_1 + h_2}{4}\right)(t_1 + t_2), \tag{6a}$$

$$\frac{\psi_1}{\psi^-} = \frac{\dfrac{h_2^*}{h^*}\left(\beta_2 - \dfrac{|\boldsymbol{K}|^2 \varepsilon}{ik}\right)}{\beta^+ - \dfrac{|\boldsymbol{K}|^2}{ik}\dfrac{\varepsilon h_2^*}{h^*}}, \tag{6b}$$

$$\frac{\psi_2}{\psi^-} = \frac{\dfrac{-h_1^*}{h^*}\beta_1}{\beta^+ - \dfrac{|\boldsymbol{K}|^2}{ik}\dfrac{\varepsilon h_2^*}{h^*}}, \tag{6c}$$

$$\frac{f w_I}{h_1^*} = -ik\beta_1\psi_1. \tag{6d}$$

We have defined β_i as $\beta - |\boldsymbol{K}|^2 U_i$ and β^+ as $\beta - |\boldsymbol{K}|^2 U^+$, where U^+ is the vertically averaged basic flow

$$\frac{(U_1 h_1^* + U_2 h_2^*)}{h_1^* + h_2^*} \quad \text{and} \quad |\boldsymbol{K}|^2 = k^2(1 + q^2).$$

Suppose that the perturbations velocities are unimportant in the lower layer temperature equations, the equilibrium between mean zonal advection and heating at the ground leads to:

$$\frac{t_2}{\mathrm{SST}_0} = \frac{\left(1 - ik\dfrac{U_2}{\lambda}\right)}{1 + k^2\dfrac{U_2^2}{\lambda^2}}.$$

This simple balance shows first that the response t_2 is always less than the excitation SST_0 whatever the value of the non-dimensional parameter kU_2/λ and second that t_2 and SST_0 are in phase when kU_2/λ is small and become in quadrature when it is large, the atmospheric response being carried downstream of the SST. The parameter kU_2/λ around O(1) for large scale waves therefore suggests a phase around $\pi/4$ between t_2 and SST_0. Although this phase prediction is of value, this simple equilibrium does not allow for any positive feedback on the SST. If any is to be found, the perturbation velocities must be large.

Eqs. (6) shows readily that this will be the case for those wavenumbers for which the denominators in (6b–6c) are small, that is for:

$$\beta^+ = 0, \quad \text{i.e., } k^2 = k^{+2} = \frac{\beta}{U^+(1 + q^2)}.$$

Because the vertically averaged velocity U^+ controls the resonant wavenumbers, this resonance is non local. Note that because the flow is sheared vertically, vertical velocities are present when this external Rossby mode is excited. One expects abrupt phase changes and large amplitudes velocities near k^+, if the relative width of the resonance curve (given by $\varepsilon h_2^*/h^*/U^+k^+$) is small a condition which hinges on the value of the bottom friction coefficient ε. Otherwise expressed as the ratio of the advective time scale at k^+ over the dissipative time scale, it is one of the central parameter of the resonant thermal response studied herein.

To compute the atmospheric response in a realistic case we have chosen parameters appropriate to the mid-latitudes (Table 1). Since this model formulation neglects zonal and meridional dependence of the mean zonal flow, some averages of that mean state must be carried out in both directions. The zonally averaged annual mean profiles which are chosen are representative of the mean state averaged around the globe between 45° lat and 55° lat (Peixoto and Oort, 1992).

The 2 levels are centered at 750 hPa and 250 hPa, respectively, the full height of the model amounting to 17.5 km. Once the mean temperatures are chosen, the densities at each level are obtained from the perfect gas law. The vertical scale of the atmosphere H_0 is chosen to be 8 km.

Given the mean state, it is apparent that this theory has two important parameters the heat exchange coefficient, and the bottom friction coefficient both of which parametrize the effects of subgrid scale motions. The heat exchange coefficient has been estimated from climatology by Haney (1971). Assuming no atmospheric change to SST changes, Haney's values were in the range of 40 W m^{-2} K^{-1}. Quite recently this estimation was revisited using 40 years of North

Table 1. *Parameter values of the two level atmospheric model*

	h_i (km)	T_i (K)	ρ_i (Kg m^{-3})	$\left.\frac{\partial \bar{T}}{\partial \sigma}+\frac{K\bar{T}}{H_0}\right\|_i$ (K m^{-1})	$\frac{\partial \bar{T}_i}{\partial y}$ (K m^{-1})	U_i (m s^{-1})
$i=1$	12.97	230	0.378	$2.94\cdot 10^{-3}$	$-8\cdot 10^{-7}$	15
$i=2$	4.6	255	1.024	$2.88\cdot 10^{-3}$	$-4.5\cdot 10^{-6}$	7.7

Atlantic observations of SST and surface heat flux by Frankignoul, Czaja and L'Heveder (1998). They showed that the heat flux sensitivity to anomalous SST varied across the domain but averaged to about 20 W m^{-2} K^{-1}, a value that is kept in the following. Parametrisation of bottom friction is carried out following standard practice of fitting an Ekman layer to a constant stress layer close to the ground allowing the bottom stress to be calculated in terms of the geostrophic velocity $\boldsymbol{u}_g$ above (Bluestein, 1992):

$$\frac{\boldsymbol{\tau}}{\rho_a} = C_D(\cos\alpha - \sin\alpha)^2 |\boldsymbol{u}_g| \cdot \boldsymbol{u}_g\, e^{i\alpha},$$

where $\alpha \sim 24^\circ$ is the turning angle estimated empirically and C_D a drag coefficient. This expression is linearized by considering a turbulent field with given characteristic velocities U_T superimposed on the mean flow:

$$\frac{\boldsymbol{\tau}}{\rho_a} = C_W \boldsymbol{u}_g\, e^{i\alpha}, \tag{7}$$

where $C_W = C_D$ $(\cos\alpha - \sin\alpha)^2 U_T$. With $C_D = 1.5\cdot 10^{-3}$ (over the ocean) and $U_T \sim 5$ m s^{-1}, $C_W \sim 1.87\cdot 10^{-3}$ m s^{-1}. By integrating the stress divergence across the lower layer and neglecting any interfacial eddy stresses, we obtain:

$$\varepsilon = \frac{C_W}{h_2} \sim 4\cdot 10^{-7}\,\mathrm{s}^{-1},$$

or an equivalent spin down time scale of 28.9 days.

Note finally that friction enters the theory only through the curl of the stress. Taking the curl of (7), we obtain C_W $(\cos\alpha\nabla\times\boldsymbol{u}_g + \sin\alpha\nabla\cdot\boldsymbol{u}_g)$. Since the flow is nearly geostrophic and α is not large we neglect the divergence term and take $\cos\alpha \sim 1$ in the expression of the curl, a procedure which amounts to neglect altogether the rotation by the angle α in (7).

Given those parameters the response of the model atmosphere to periodic SST can be estimated. An aspect ratio ($q=2$) has been chosen for meridional/zonal wavenumbers. Although the wavenumbers are integral multiples of $1/R_T \cos\theta_0$, we present a continuous distribution to show the details around the resonance. We first show the wavenumber distribution of the gain, the real part of $[t_2/\mathrm{SST}_0]$ in Fig. 2. This quantity measures that part of the response that is in phase with the SST and is obviously the relevant quantity to discuss thermal feedbacks. At wavenumber smaller than 1 the gain is near unity, the advection terms are negligible, the atmospheric temperatures nearly equal to SST. As the zonal wavenumber increases, the gain slowly decreases and this is the regime (described previously) of balance between zonal advection and air–sea heat exchange. As the wavenumber approaches k^+, near resonance conditions

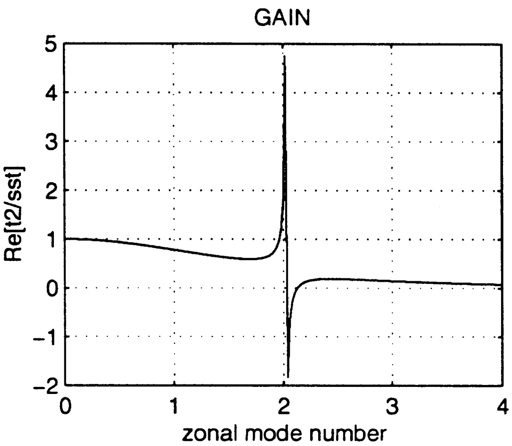


Fig. 2. The ratio of the lower level atmospheric response to prescribed SSTs is shown as a function of the zonal mode number of the SST distribution. The gain is less than 1 for nearly all mode numbers but for a very narrow window on the low wave number side of external Rossby wave resonance. In this window the atmosphere acts as an amplifier. The resonance is sharp because the perturbation advective time scale is much less than the spindown time due to bottom drag.

are found and the gain increases to about 5. It decreases as rapidly on the high wavenumber side and remains thereafter to a small value. What is most striking in this figure is the sharpness of the resonance curve because it provides a powerful wavenumber selection. The observed relative width in wavenumber $\Delta k/k$ corresponds to the theoretical width estimated previously

$$\frac{\varepsilon}{U^+k^+}\frac{h_2^*}{h^*} \sim 3.45 \cdot 10^{-2}.$$

It is small because the advective time scale (around 2 days) is at least an order of magnitude less than the bottom friction time scale so that the mode is nearly inviscid. To understand physically how amplification appears around $m = 2$, the phases of ψ_2 and t_2 perturbations are shown in Fig. 3. Note that ψ_2 and t_2 are in opposition of phase for nearly all wavenumbers but those near resonance. It is only in a very narrow window of wavenumber space that high (low) pressure are to be associated with warm (cold) anomalies at the lower level. Because of thermal wind, the upper layer pressure anomalies (in phase with lower layer) are larger and the perturbations are equivalent barotropic. Previous investigation of mid-latitude ocean atmosphere coupling that assumed proportionality between surface pressure and temperature have therefore implicitly assumed a resonance condition to be valid for all wavenumbers, preventing the wavenumber selection associated with near resonance to occur. Looking at the balance of terms in the temperature equations at $m = 2$ (Fig. 4), shows that it is the meridional eddy velocities acting on the mean meridional temperature field that are responsible for enhancing the lower layer temperature anomaly (the term $-v_2(\overline{\mathrm{d}T_2}/\mathrm{d}y)$ is large and in phase with t_2). The bottom heating is in opposition of phase with the atmospheric temperature while the zonal and vertical advection terms are nearly in quadrature. The temperature anomalies are small at the upper level because even though the meridional velocities are large, the mean meridional temperature gradient is negligible there. What is left of that meridional advection is balanced by vertical advection. Away from resonance at either lower or higher wavenumbers, the lower level temperature field has similar phase relationships with the SST (around $\frac{1}{4}\pi$ displacement downstream) but the pressure is now shifted between $[\frac{1}{4}\pi, \frac{1}{2}\pi]$ upstream, so that meridional temperature advection actually is a negative feedback for the temperature anomaly. In the absence of friction, (6c) shows that ψ_2 and ψ^- (that is $t_1 + t_2$) have the same phase (relationship) in a wavenumber interval:

$$k_1 < k < k^+.$$

Since

$$k_1 = \left(\frac{\beta}{U_1(1+q^2)}\right)^{1/2}$$

is smaller than k^+ (because the upper layer velocities U_1 are larger than U^+), the right phase relationship for amplification (i.e., equivalent barotropic structure) is in a slot on the low wavenumber side of the resonant wavenumber k^+ (Fig. 3d). This is not a sufficient condition for amplifications though, and gain greater than 1 occurs only for large enough meridional velocities with k close to k^+. One further remark concern vertical velocities, which are far from negligible. Their roles in the vorticity balance are central. At wavenumbers less than k^+, the balance becomes more Sverdrup-like with advection of planetary vorticity equilibrating stretching while at wavenumbers higher than k^+, the β effect becomes less and less important leaving a balance between advection of relative vorticity and stretching. The terms of the vorticity equations are shown in physical space for $m = 2$ in Fig. 5. At both levels, the zonal advection of relative vorticities contributes much to the vorticity tendency. It is nearly in phase and of the same sign as the stretching term in the lower layer, these two contributors balancing a large advection of planetary vorticity, a clear departure from Sverdrup balance. If we consider the situation relative to the SST maximum (with geopotential maximum some distance downstream), vertical velocity is positive, stretching the lower layer upstream of the SST maximum in a region of poleward velocities, the sum being associated with a vorticity deficit. The deficit is large because it has to compensate for positive gain of vorticity through the advection by the mean zonal flow.

Although this near resonance process appears to be a powerful wavenumber selection for amplification of the response to periodic SST perturbation, we must test its robustness to variations of parameters. There are two important parameters, which concern us here, the bottom friction ε, and the heat exchange coefficient γ. The gain of the

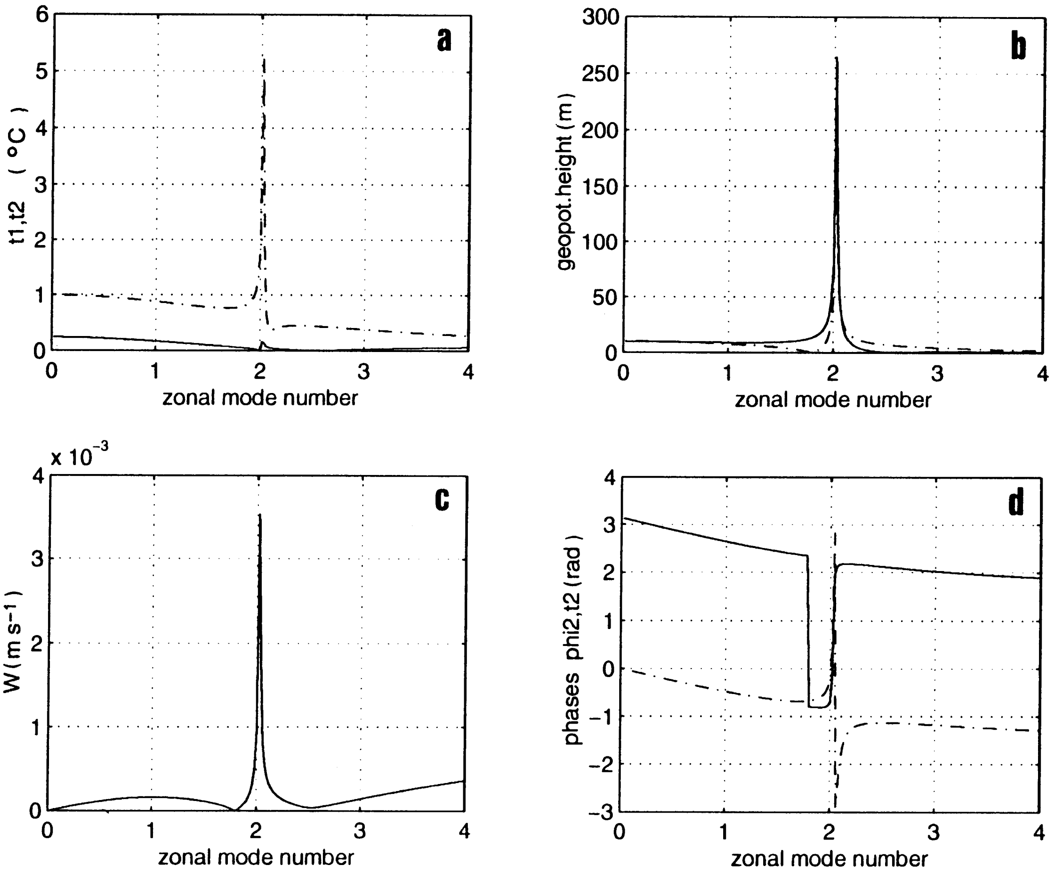


Fig. 3. Description of the response to a 1°C SST perturbation as a function of mode number. Temperature (a) and geopotential heights (b) at respectively upper levels (continuous), lower level (dashed). (c) Vertical velocity at the interface. (d) Phase differences between the lower layer geopotential (continuous) and temperature (dashed). Only in a very narrow window is the atmospheric response equivalent barotropic. Elsewhere the response is baroclinic associating positive temperature anomalies and low geopotential.

system has been computed at $m = 2$ as a function of γ and the effective drag coefficient C_W in Fig. 6. The gain is seen to be greater than 1 in a significant region of parameters space of geophysical interest. It shows first of all that an inviscid theory does not lead to amplification and that some amount of surface drag is necessary. In the whole range of realistic values of C_W [10^{-3}–$6 \cdot 10^{-3}$ m s^{-1}], amplification occurs provided that γ is large enough [>5 W m^{-2} K^{-1}]. Optimum gain is found for values of $C_W \sim 1.5 \cdot 10^{-3}$ and $\gamma \sim 10$ W m^{-2} K^{-1} which are on the low side of the range of geophysical interest. The amplification factor decreases for γ and ε in excess of these optimal values. For a given value of γ a high drag coefficient broadens the resonance so much as to make the response less than the excitation. For a given value of the drag coefficient, the gain decreases for γ beyond the value for the optimum gain. In such a case the SST and the response t_2 become more in phase and as a result the phase lag between the lower layer geopotential and the SST is reduced. Then the meridional geostrophic velocities become close to quadrature with the SST and the main amplification factor for the lower layer temperature perturbation disappears. This sensitivity study shows unambiguously that near resonant amplification occurs for all reasonable values of the heat exchange and bottom drag coefficient. On the other hand, the gain varies

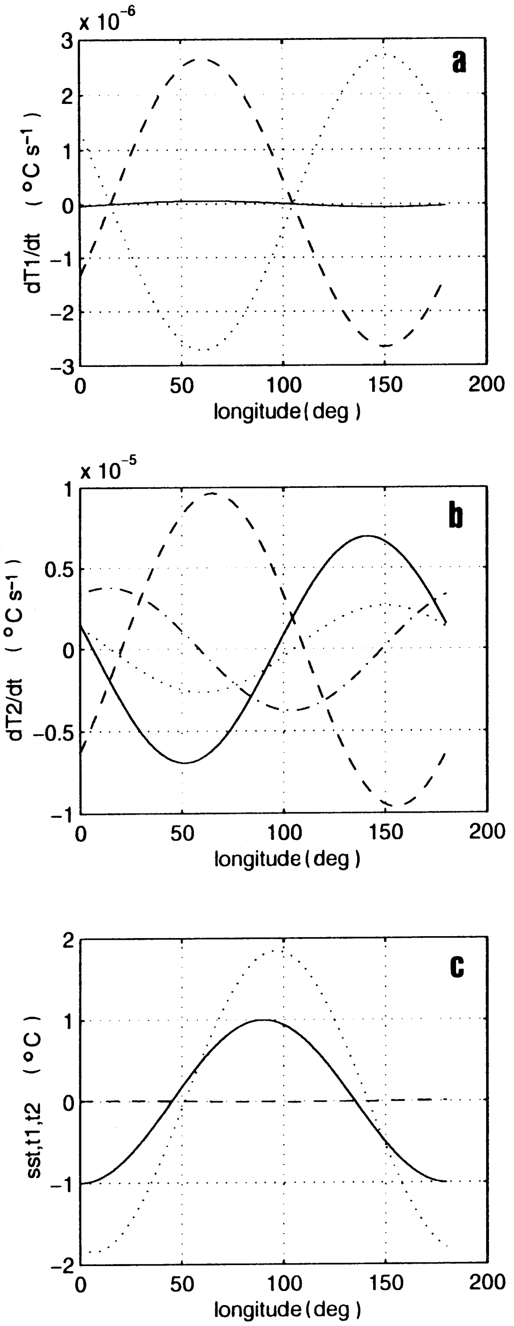
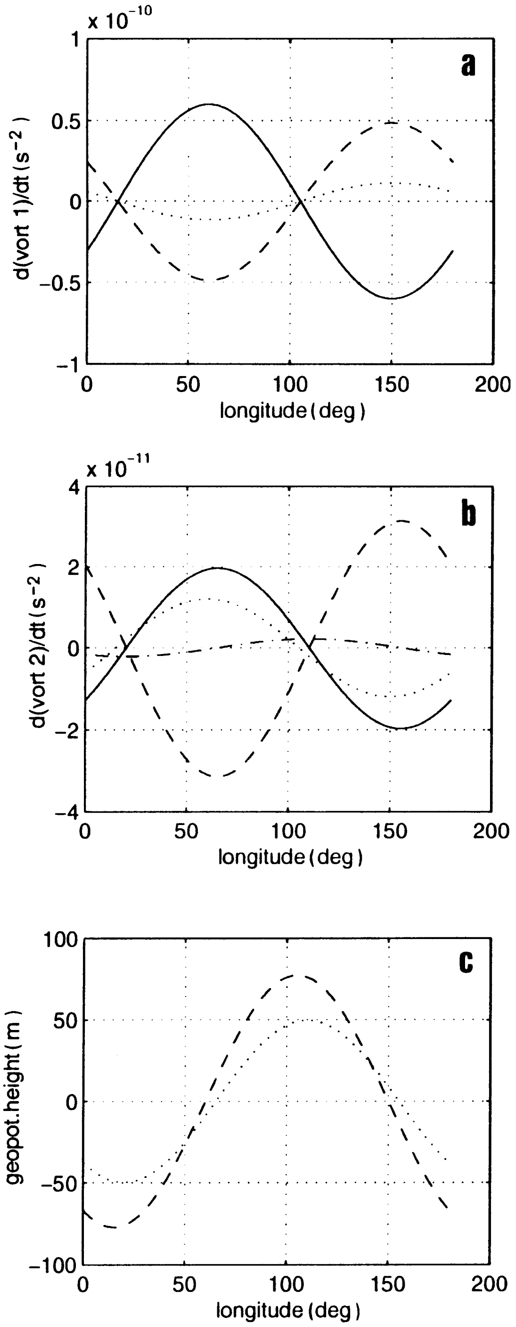


Fig. 4. Heat balance of terms for the zonal mode number 2 written as contributions to temperature temporal changes for the upper (a) and lower level (b). $-U_i(\mathrm{d}t_i/\mathrm{d}x)$ (continuous), $-v_i(\mathrm{d}\overline{T_i}/\mathrm{d}y)$ (long dash), $-w_1(\mathrm{d}\overline{T}/\mathrm{d}\sigma + K\overline{T}/H_0)$ (short dash), thermal forcing (—·). The SST (continuous), the lower (short dash) and upper level temperatures (long dash) are shown in (c). Note how the strong advection by poleward perturbation velocities reinforce the positive SST anomaly. This is a nearly equivalent barotropic situation with high geopotential at both levels downstream of the positive SST anomaly.

widely with uncertainties in these coefficients. Roads's (1982) remarks that "thermal damping in combination with surface friction can actually increase the amplitudes near resonance" point to a common physical explanation. The origin of this surprising amplification that requires friction can be explained with the energetics. Multiplying the perturbation temperature eq. (5c) by temperature t_i, summing over the layers and zonally averaging shows a three term balance of the available potential energy (APE) equation for the perturbations. There is first a conversion term between APE and kinetic energy, which is negative (dissipative for APE) if dissipation in the vorticity equation is present. The second is the forcing term $\gamma\langle t_2(\mathrm{SST} - t_2\rangle$ which is also negative if the gain is greater than one and the third is the familiar $-\langle v_i t_i\rangle(\mathrm{d}\overline{T_i}/\mathrm{d}y)$ dominated by values in the lower layer. For this last term to be positive and equilibrate the other two, the perturbation heat fluxes must be "downgradient". Only in this situation can enough mean APE (associated geostrophically with the mean zonal flow) be extracted to force a strong planetary wave response. In the absence of vorticity dissipation, eqs. (6b) and (6c) show that the stream functions at upper and lower levels are in phase with the temperature fields so that the perturbation heat fluxes vanish. With some dissipation, phase lags appear allowing release of the mean APE. This explains the vanishing gain observed in Fig. 6 as the bottom drag goes to zero.

This simple physical analog has response characteristics which may perhaps help to interpret physically the rich responses of atmospheric GCM. Let us summarize them:

(1) Significant ("equivalent barotropic") response is expected if the scale of the prescribed SST is in the wavenumber window $[k_1, k^+]$ mentioned previously. In case of periodic SST distribu-



tion, the wavenumber $m/R_T \cos\theta$ associated with an integer mode number m must be in that interval. The respective values of $U_1(U_2)$ of $15\ \mathrm{m\,s^{-1}}$ ($7.7\ \mathrm{m\,s^{-1}}$) used for the previous figures guaranteed that the mode number 2 was in the window [1.8, 2.06]. This condition of course depends on the vertical structure of the mean zonal flow because the wavenumber window narrows when the bottom velocity increases towards the value of the upper layer flow. Control of the vertical structure of the mean flow by seasonal effects and the turbulent transient eddy cascade are then expected to bear upon the planetary wave response.

(2) For forcing outside the wavenumber window, the response is baroclinic (low pressure downstream of warm/SST). This essentially forced response is expected to decrease as wavenumber increases.

(3) The amplitudes of the response near the resonance are controlled by the value of the surface drag. A large sensitivity to dissipation values of vorticity is to be expected and it is only in a small range of dissipation values that significant release of mean APE (leading to large responses) is found.

In the off-resonance case turbulent dissipation can be neglected and at exact resonance the large scale dynamics cancel out. The near resonant case that we focus on in the wavenumber window $[k_1, k^+]$ was more difficult to predict because no terms could be neglected a priori. Finally because of the nature of the near resonant response proposed here, initial value problems predict algebraic growth at most linear in time of the response amplitudes, rendering the identification difficult. The situation will be vastly different in a coupled situation: indeed if the process survives coupling with an active ocean, exponential growth is to be expected.

Fig. 5. Vorticity balance of terms for the zonal mode number 2 is written as contributions to vorticity temporal changes for the upper (a) and lower level (b). $-U_i(\partial/\partial x)\xi_i$ (continuous), $-\beta v_i$ (long dash), stretching (short dash), bottom drag (—·). Geopotential height streamfunction for the upper (long dash) lower level (short dash) are shown in (c). Note how the advection of perturbation vorticity by the mean zonal flow perturbs the Sverdrup equilibrium when the atmospheric response is in the equivalent barotropic window.

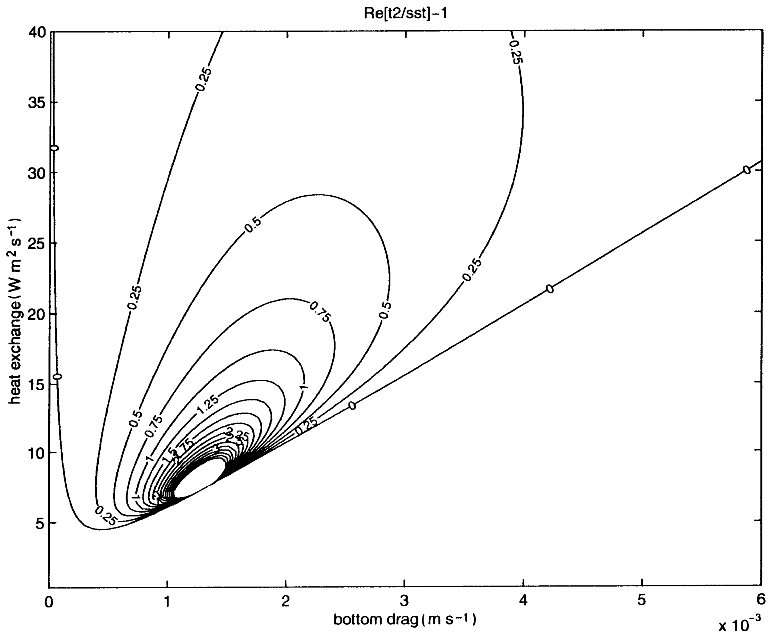


Fig. 6. Sensitivity of the gain to two external parameters, the linearized bottom drag coefficient (in m s^{-1}) and the heat exchange coefficient. When the plotted quantity (= gain − 1) is positive, the atmosphere act as an amplifier. Both external coefficients must be different from 0 for amplification showing the control of the large scale response by the small scale turbulence. If they are too large (beyond values of geophysical interest) the amplification weakens and even vanishes (large bottom drag).

3. Coupling to an oceanic model

As mentioned earlier the amplification factor for temperature found at resonance turns out to be a positive feedback for the SST in the oceanic compartment. Furthermore the anomalous winds drive anomalous oceanic velocities that may modify or annihilate this resonance feedback. To find out the robustness of the process, we need to couple the atmospheric model to an ocean model. To keep matters as simple as possible, a zonally periodic channel geometry is chosen for the ocean, so that the application of such single Fourier component analysis is restricted to regions of the Antarctic circumpolar current (ACC) system. This strong eastward current, which carries more than 150 Sv around the globe, is of course the major feature to take into account. Vertically sheared, it is unstable and generates a field of active mesoscale eddies. At the planetary wave scale of interest here, oceanic advection of relative vorticity can be safely neglected because the oceanic velocities are at least one order of magnitude less than atmospheric velocities. The so called planetary geostrophic approximation widely used in oceanic modeling makes use of this fact and of the small internal Rossby radius: large scale barotropic Rossby waves (that may be stationary in the atmosphere) propagate at phase speed much larger than oceanic speed and therefore cannot be resonant in the ocean. However the essentially non dispersive baroclinic Rossby waves which have phase speed in the range of the large scale ocean circulation O (1 cm s^{-1}) can be made stationary as shown by Schopp (1991). We believe that this interesting possibility does not occur in mean flows such as the ACC where the mean velocity (meridionally averaged) is much larger (about 10 cm s^{-1}). In this regime, the large scale oceanic perturbations on a mean zonal flow decreasing with depth are top-trapped modes that look somewhat like the f plane boundary waves that move with the speed of the mean flow at their penetration depth (see Gill, 1982, p. 550). Because of the near surface intensification of these modes, it makes sense to choose for the ocean a reduced gravity model with only one active layer whose depth becomes, of course, an additional external parameter.

Because of later coupling to the atmospheric 2-level model, the oceanic variables are labelled with index 3. The same vertical discretization presented earlier for level 1, 2 is also used for level 3 (Fig. 1). It is assumed that the oceanic mean state is a zonal geostrophic jet, which varies in the vertical only. Similar discretization of the hydrostatic relation (Section 2) is carried out to ensure energy conservation for the two levels, 3 and 4, the deepest of which will be put to rest for the perturbations:

$$P_3 - P_4 = \rho_0 \alpha g \frac{h_{0C}}{2} T_I, \qquad (8)$$

with

$$T_I = \tfrac{1}{2}(T_3 + T_4).$$

h_{0C} is the sum of the layer depths $h_3 + h_4$ and α the thermal expansion coefficient for water (omitting salinity effects consistently with the neglect of moisture transports in the atmospheric model). The net momentum of the upper layer is now forced by the stress at the air–sea interface:

$$f\boldsymbol{k} \times \boldsymbol{u}_3 = \frac{-\nabla P_3}{\rho_0} + \frac{\boldsymbol{\tau}}{\rho_0 h_3}, \qquad (9)$$

whose curl is:

$$\beta v_3 = \frac{-f w_I}{h_3} + \frac{\nabla \times \boldsymbol{\tau}}{\rho_0 h_3}, \qquad (10)$$

w_I being the vertical velocity at the base of layer 3. Finally the perturbation temperature equation linearized around the mean geostrophic zonal flows U_3 is:

$$\frac{\partial t_3}{\partial t} + U_3 \frac{\partial t_3}{\partial x} + v_3 \frac{\mathrm{d}\overline{T_3}}{\mathrm{d}y} + w_I \frac{\Delta T_0}{h_3} = \lambda'(T_{\mathrm{atm}} - t_3), \qquad (11)$$

where

$$\Delta T_0 = \tfrac{1}{2}(\overline{T_3} - \overline{T_4}) \quad \text{and} \quad \lambda' = \gamma/\rho_{0C} C_p^0 h_3.$$

Eqs. (8) through (11) can be cast into an equation involving only the temperature variable after using (9) and (10) to express v_3, w_I in term of p_3, itself related to t_3 by (8):

$$\frac{\partial t_3}{\partial t} + (\tilde{U} + C_R) \frac{\partial t_3}{\partial x} + w_E \frac{\Delta T_0}{h_3} + v_E \left(-\beta \frac{\Delta T_0}{f_0} + \frac{\mathrm{d}\overline{T_3}}{\mathrm{d}y} \right) = \lambda'(T_{\mathrm{atm}} - t_3), \qquad (12)$$

where:

$$\tilde{U} = U_3 + \frac{g\alpha h_{0C}}{4f_0} \frac{\mathrm{d}\overline{T_3}}{\mathrm{d}y},$$

$$C_R = -\beta \frac{g\alpha \Delta T_0}{4f_0^2} h_{0C},$$

$$w_E = \frac{\nabla \times \boldsymbol{\tau}}{\rho_0 f_0}, \qquad v_E = \frac{-\tau_x}{\rho_0 h_3 f_0}.$$

The temperature evolution given by (12) for the upper oceanic layer is governed by processes that are all well-known.

(1) Zonal advection by the velocity $\tilde{U}$. This is not the mean speed of the upper layer but one more closely related to some average of layer 3 and 4, mimicking the steering level in a reduced gravity discretization. In the case of eastward (westward) advection, the westward phase velocity C_R of the non dispersive Rossby waves decrease (increase) the net advection.

(2) Vertical heat advection by Ekman pumping acting across the interface at the base of layer 3.

(3) Meridional advection by Ekman velocities (at right angle to the wind stress). Note that the mean temperature gradient is corrected by a term function of β. This term can alternatively be combined with the previous one after redefining Ekman Pumping as $\nabla \times (\boldsymbol{\tau}/f)$.

(4) Heat exchange at the air–sea interface.

Besides the advection terms which simply transport temperature anomalies at the speed $\tilde{U} + C_R$, this simple oceanic model identifies two interactions terms which involve the wind stress and one which involve the heat exchange at the interface. In order to find out their active or passive role in driving the oceanic anomalies, the oceanic model needs to be coupled to the atmospheric model. The major coupling hypothesis is that since the atmosphere adjust rapidly to oceanic temperature anomalies, the time dependent oceanic model (12) can be coupled to the steady state atmospheric model (5). A measure of this approximation is given by the three orders of magnitude that separate the oceanic thermal inertia from that of the atmosphere as shown by the ratio

$$\frac{\lambda'}{\lambda}\left(=\frac{\rho_a C_p^a h_2}{\rho_0 C_p^0 h_3}\right).$$

This low frequency coupling approximation guarantees that the instability that will eventually be found is a truly coupled instability and differs from internal instabilities of the atmosphere which are filtered out. Furthermore because the atmosphere adjusts instantaneously to the changing lower boundary condition, atmospheric perturbations are then advected along by oceanic anomalies. Consistent with our formulation that integrates vertically across boundary layers on each side of the air–sea interface, the coupled model governed by (5) and (12) is "closed" by extrapolating mid-level temperature to the air–sea interface as:

$$\mathrm{SST} = t_3, \qquad \text{in (5c)},$$

$$T_{\mathrm{atm}} = t_2, \qquad \text{in (12)}.$$

The stress in (12) which is continuous at the air–sea interface is still given by (7) (with rotation neglected for reasons given previously) so that the conservation of momentum holds. After the same Fourier decomposition of t_3 as $\Re[e^{st} e^{ik(x+qy)}]$, we satisfy (5) and (12) by expressing t_2 and τ in terms of t_3 from the atmospheric model equations. After eliminating t_2 and τ in (12) one obtains a single linear equation for t_3 whose coefficient must vanish yielding the eigenvalue relation governing s. When wind stress effects are ignored, it is immediate to show that the real part s_r is due to the heating term in (12), $\lambda'\Re(t_2 - t_3)$, so that in the resonance window mentioned previously (Section 2) where the gain is larger than one, s_r is positive. We may also expect that the zonal advection term in (12) provides the essential contribution to the imaginary part s_i:

$$s_i = k(\tilde{U} + C_R).$$

In this simple limit, the transport of the anomaly is carried out by the ocean while the amplification is carried out through the resonant thermal response of the atmosphere to the oceanic anomaly described previously. Exponential amplification through atmospheric resonance favors a given wavenumber, while the period is essentially set by passive oceanic advection. We now discuss if this limit is of some value for the existence of the Antarctic Circumpolar Wave of White and Peterson (1996).

4. Applications to the Southern Ocean

In the Southern Ocean, White and Peterson (1996) have identified over a 15-year period a remarkable zonally periodic SST wave train that propagates along 56°S at an average eastward velocity of 6–8 cm s^{-1}. Altimetric heights in phase with the SST have confirmed the existence of this extraordinary mode (Jacobs and Mitchell, 1996). This is essentially a mode number 2 whose period is about 4–5 years. Travelling along with the SST anomalies (0.5°C) are SLP (sea level pressure) and MWS (meridional wind stress) anomalies with respective amplitudes of 4 hPa and 0.15 dyn cm^{-2}. Warm anomalies lead high SLP by about 90°. The anomalies have a meridional/zonal wavenumber aspect ratio around 2. The horizontal pattern shows warm SST anomalies to be sheared in a spiral towards the equator. The belt of maximum pressure anomalies is observed to be shifted somewhat polewards of the SST anomalies. Such a coherent structure needs amplification to sustain itself against eddy diffusion and the process of thermal resonance described earlier is shown to provide amplification within a realistic range of external parameters of the model (surface drag, heat exchange coefficient, mean states).

Before testing this idea, however, we asked whether a large scale internal instability of the Antarctic circumpolar current itself could not be the cause of the ACW. Huck et al. (1999) and Colin de Verdière and Huck (1999) have shown that decadal variability could appear in idealized ocean models through large scale baroclinic instability localized in the North west corner of their rectangular oceanic basins. To investigate this possibility a linear planetary geostrophic instability calculation was carried out using a zonal eastward jet (varying only in the vertical) and stratification appropriate to the ACC. Internal instabilities were ruled out because at no value of the dissipation could any significant growth rate be detected at planetary scales.

Since the model that is now proposed only includes zonal dependence, some meridional averaging of the oceanic mean state is appropriate. Vertical profiles of zonal velocity and Brunt

Väisälä frequency averaged over the [40°S–60°S] latitude band give a good representation of the ACC along the Greenwich meridian. This is the Ajax section described by Whitworth III and Nowlin (1987). The vertical intensification of the mean flow provides the surface intensification of the linear wave modes. We envision the SST anomalies of the ACW as the surface expression of one such mode. Unfortunately there is no evidence of the value of vertical penetration in the observations and so the depth h_3 of the active layer of the ocean model is therefore an additional parameter. The effects of varying h_3 are straightforward however: reducing h_3 essentially increase the oceanic forcing through increase of the divergence of heat fluxes and wind stresses. It also modifies the advective velocity. For a chosen value of h_3, U_3 was estimated by simply averaging of the current velocity over that depth while the Brunt–Väisälä frequency was averaged around the base of that interface. Because the ACC includes also a salinity front that opposes the temperature front, the vertical shear of the zonal current was used to provide an equivalent geostrophic temperature gradient using a standard thermal expansion coefficient. Note furthermore that a depth h_{OC} sum of layer thickness 3 and 4 occur in the hydrostatic relation (8). Although the perturbations are neglected in layer 4, the mean flow is still present. We used the observed mean flow depth scale distribution to calibrate the hydrostatic relation between layer 3 and 4. The parameters that describe the ocean model are summarized in Table 2. The parameters of the atmospheric model are as in Table 1. The lower layer velocity U_2 has been chosen so that the mode number 2 falls in the wavenumber window appropriate for equivalent barotropic response.

Using these parameters, the velocity C_3 ($= \tilde{U} + c_S$) is equal to 6.2 cm s^{-1}. It is the eastward phase velocity of free non-dispersive planetary waves confined in layer 3. However through thermal and wind stress action, forced waves may have a phase velocity different from C_3. The eigenvalue relation obtained from eqs. (5) and (12) has a single root since there is only one prognostic equation, namely the oceanic temperature equation. It is numerically evaluated and illustrated in Fig. 7. The different processes contributing to the real part of s are shown as a function of the zonal mode number m in Fig. 7a. It is readily seen that only the mode number 2 is unstable and that the thermal forcing dominates, reminiscent of the steady state atmospheric response to SST discussed in Section 2. To confirm this, the other contributors to the growth rate in eq. (12) the Ekman Pumping and the Ekman velocity* advection term are also shown in Fig. 7b. For all wavenumbers considered, the thermal forcing dominates by more than one order of magnitude over the wind stress effects. Near $m = 2$ (in the bandwidth noted in Section 2) the growth rate is entirely due to the thermally resonant response. Both wind stress effects are positive feedbacks there but quite weak compared to the thermal forcing. The imaginary part of the eigenvalue relation shown in Figs. 7c, d, indicates that the thermal forcing modifies also the unforced value C_3 of the eastward velocity, the mechanical forcing remaining unimportant. The coupling increases the eastward velocity of the coupled wave above that of the oceanic free wave but note the exception where the coupling slows the phase speed in the vicinity of the resonance. The slowing of eastward propagation is caused by non negligible perturbation velocities for it is simple to show that zonal advection with thermal coupling alone leads to enhanced eastward phase speed ($> C_3$) in westerly flows. This "projection" of the thermal forcing on the speed of SST anomalies was noted by Frankignoul (1985) using a 2-level atmospheric model similar to that of Egger (1977). The linear coupled model does not require much tuning to reproduce the observations. Insisting to have positive growth rates for integer wavenumber and zonal speed of the anomalies in the range of observed ones determines a narrow range of possible low level wind U_2 (all other parameters being

Table 2. *Parameter values of the oceanic model*

h_3 (m)	h_{OC} (m)	U_3 (cm s^{-1})	$\frac{d\bar{T}_3}{dy}$ (°C m^{-1})	ΔT_0 (°C)
200	1000	7.2	$-2.1 \cdot 10^{-6}$	1.6

* The meridional Ekman velocities have been computed by choosing the "equatorward" zonal branch of the atmospheric pressure field consistent with the White et al. (1998) observations that the atmospheric pressure field has also a poleward shift relative to the SST field.

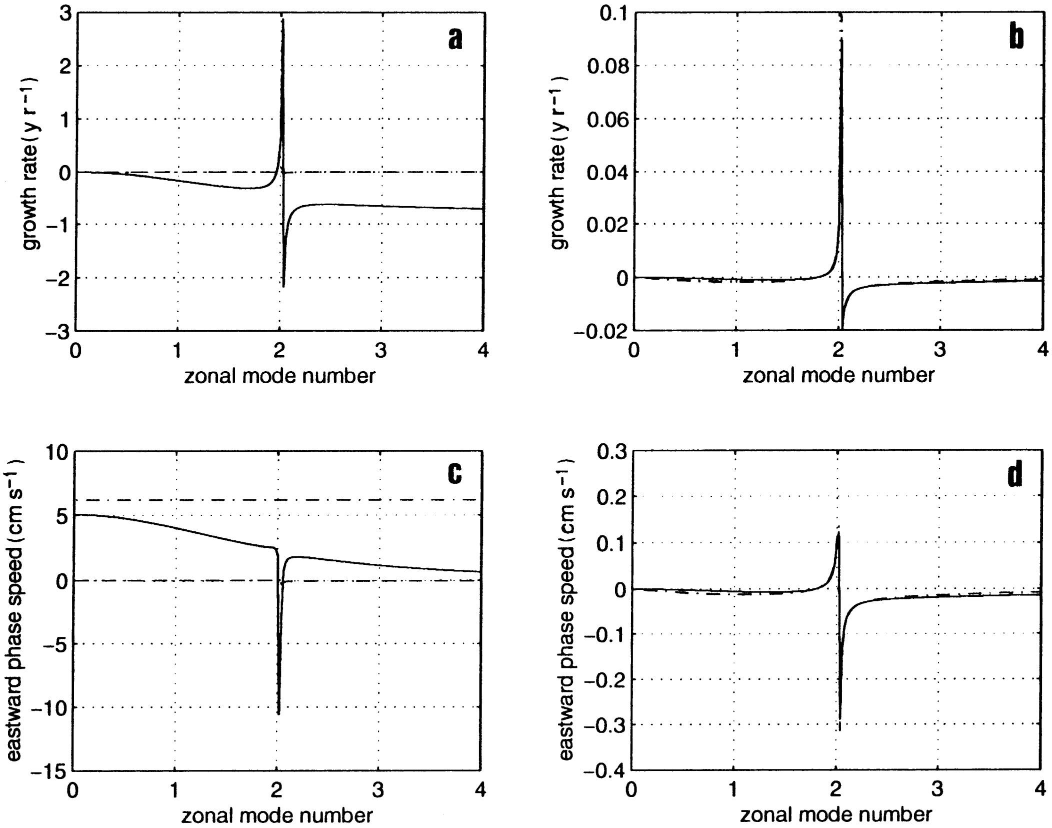


Fig. 7. As a function of zonal mode number, the figure shows: (a) the various contributors to the growth rate with the thermal forcing (continuous) dominating over wind stresses (dashed); (b) the individual smaller contributions to the growth rate of the wind stress curl (continuous) and Ekman velocities (dashed); (c) the various contributions to the eastward phase speed of the anomaly: overall advection by the ocean (—·), thermal forcing (continuous) and smaller wind effects (dash); the thermal forcing increases the eastward phase speed away from resonance while the opposite happens near resonance. Because of this sensitivity, the precise phase speed is difficult to predict but we have always found an increase of phase speed caused by the forcing in unstable situations; (d) the individual smaller contributions to the phase speed of the wind stress curl (continuous) and Ekman velocities (—·).

kept constant). With $U_2 = 7.7$ m s^{-1}, the eastward phase velocity is 8.4 cm s^{-1} (this is within the regime where the coupling increases the phase velocity in Fig. 7c) giving a period of 4.8 yr and a growth rate of 0.65 cpy for mode number 2.

Given the uncertainties in model parameters, it is important to evaluate the robustness of the coupled process through a sensitivity analysis. Fig. 8 shows the growth rates with respect to variations of the lower layer wind velocity and the bottom drag coefficient, keeping constant the mean temperature field and the mode number ($m = 2$). Note the large region of positive values which is adjacent to a region of equally high negative values when the other side of the resonance is reached. As a consequence of this dispersive property of the growth rates, it is difficult to predict what a net growth rate would be in a geophysical situation when the mean state changes seasonally. We can, however, verify that the result is robust to additional causes of eddy damping. The dashed curve in Fig. 8 delineates the region of positive growth rates when the heat diffusivities by the transient eddies in the atmosphere

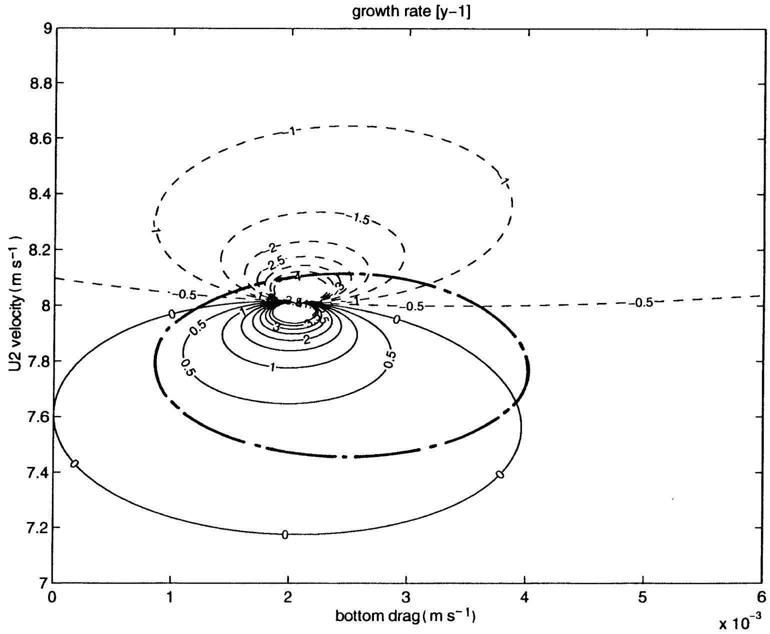


Fig. 8. The growth rate as a function of the lower level wind velocity and bottom drag coefficient. When the lower wind varies, the upper wind is changed according to the thermal wind relation with fixed temperature gradients. The mode number is fixed ($m = 2$) and all other parameters are defined in Table 1 and in the text. Note the sharpness of the resonance and the large growth rates that result. The domain inside the thick dashed curve is the reduced domain of instability when synoptic eddy diffusion of heat is added in the atmosphere ($5.5 \cdot 10^5$ m^2 s^{-1}) and in the ocean (10^3 m^2 s^{-1}).

(5.5×10^5 m^2 s^{-1}) and ocean (10^3 m^2 s^{-1}) are included. Although such parameterizations are crude the large scale of the unstable waves shows that the process isolated here survives to interactions with eddy fields producing O (1 cpy) growth rates.

This thermal near resonance theory is rather different from others presented so far that rely instead on wind induced anomalous heat transport by ocean currents, Qiu and Jin (1997), Talley (1999), Goodman and Marshall (1999). That these wind induced feedbacks are positive is no surprise. In the $[k_1, k^+]$ wave number window, high pressure occurs some distance downstream from warm SST. Anomalous winds then induce negative Ekman Pumping (ie. positive SST tendencies) at the position of the pressure maximum (The meridional Ekman transports associated with the "equatorward" branch of the anticyclone acting on the mean meridional temperature gradient produce similar effect). These translate to positive feedbacks when the phase lag between SST and atmospheric pressure is less than $\frac{1}{2}\pi$. The thermal resonance process requires similarly that the phase lag between SST and lower air temperature be less than $\frac{1}{2}\pi$. These phase requirements are somewhat in contradiction with White and Peterson's observations, who usually indicate quadrature between SST and SLP. Either process, thermal or mechanical, would predict that the ACW is then in a near neutral state. The present linear theory falls short of predicting the amplitudes of the ACW and an immediate objective is to elucidate how such patterns stabilize at finite amplitudes by keeping track of the evolution of the mean flow shown in Chen and Trenberth (1985) to be strongly modified near resonance.

The dominance of thermal over mechanical forcing can be appreciated from the perspective of a scaling argument. Let us scale the same way the temperature anomaly at low levels in the atmosphere and in the ocean. The ratio of the thermal forcing to the Ekman pumping heat advection in (12) becomes then the product of 3 adimensional numbers:

$$\frac{f_0^2 L^2}{N^2 h_{0C}^2} \times \frac{g\alpha h_{0C}}{R} \times \frac{\gamma}{C_W \rho_A C_P^0}.$$

With $L \sim 2000$ km the length scale of the anomaly and h_{0C} the scale depth of the oceanic mean flow, the respective mid-latitude estimate of the 3 factors $(4.10^{+3}, 10^{-2}, 2)$ shows readily that thermal forcing dominance is associated with small values of the oceanic Burger number (the inverse of the first factor). The near resonant atmospheric Rossby waves have a scale that is two orders of magnitude larger than the oceanic internal Rossby radius of deformation, making the Burger number very small. There is not much freedom with the two other factors at the exception of the drag coefficient C_W. At the very least this crude argument that neglects phase information, cautions against the importance of mechanical feedbacks.

It should be possible to judge the respective merits of existing theories in coupled GCMs or observations based on the following energy criteria. In our thermal near resonance process, the wave heat fluxes are down gradient in the atmosphere, releasing the mean APE associated with the jet stream whereas in the cases discussed by Qiu and Jin (1997) and Goodmann and Marshall, (1999) the wave heat fluxes release oceanic APE associated with the presence of the thermocline. Southern hemisphere atmospheric standing waves were shown by Van Loon et al. (1973) to be mostly in phase in the vertical (in comparison with the Northern hemisphere). We believe that this does not refute the present theory because a standing wave analysis is not appropriate and also because the analysis should be carried out with data exhibiting the ACW in the first place! Note finally that the vertical phase lag of atmospheric perturbations induced by surface drag is a very small signal to detect (Fig. 5c). The situation is likely to be even more difficult in the ocean and it will be some time before large scale observations of the vertical structure of the perturbations of the Antarctic circumpolar current will allow to reveal the existence of vertical phase lags. More limited prospects are simply to compute the large scale wave fluxes in a coupled GCM which simulates the ACW, to find out in which fluid are their directions down gradient.

5. Summary

We have shown that a theory of the linear response of the atmosphere to planetary scale SST distributions may be of value to interpret the diverse GCM responses that have appeared in the literature. Within the limitations of a two level vertical discretization, the response is usually baroclinic, the equivalent barotropic response being restricted to a small wavenumber window on the low wavenumber side of the resonance that occurs when a train of atmospheric Rossby waves is made stationary against the vertically averaged mean zonal wind. Near resonance conditions are the most difficult to analyze because all terms in the vorticity equations are potentially important. Sensitivity analysis of the near resonant response shows that amplification of the SST field can occur for a realistic range of heat exchange coefficient and turbulent drag coefficient. It does not exist if dissipation is absent and as such differs from previous works. This amplification favors the equivalent barotropic response that has been found in high resolution GCMs. As an example Peng et al. (1995) obtained two very different responses (respectively equivalent barotropic and baroclinic) to North Atlantic SST in a November and January initialization of their atmospheric model. They came to the conclusion that the atmosphere over the SST had warmed considerably in the November conditions providing positive feedbacks for the SST. Qualitatively at least the present theory would suggest that the scale of their SST anomaly was in the relevant "near resonance" window in the November case but not in the January case.

After coupling this 2-level atmosphere to a one level ocean, we have identified an unstable coupled wave which occurs at wavenumbers which are nearly resonant for external atmospheric Rossby waves. Thermal feedbacks play a central role to explain the existence of this air–sea coupled unstable wave whose structure is compatible with what is known about the ACW and are robust against smaller scale damping mechanisms. Both thermal feedbacks and mechanical feedbacks through heat transport by wind induced currents are positive, with the former dominating by more than one order of magnitude. Scale analysis suggest that this thermal dominance is a consequence of the low value of the oceanic Burger number, the scale of the unstable coupled anomalies being two orders of magnitudes larger than a typical oceanic internal Rossby radius. The ocean provides essentially a mechanism for moving the anomaly around the globe and setting the period. The phase velocity is modified to some extent by the coupling which speeds up the waves except in a narrow interval near the resonance. It is possible that a knowledge of the penetration depth of the ACW into the ocean (unknown so far) would help disentangle advective effects from forcing effects in establishing the period. The source of energy of the coupled instability is the atmospheric APE associated with the westerly jet. To compare this theory with two previous ones by Qiu and Jin (1997) and Goodmann and Marshall (1999) who identify the oceanic APE as the primary energy source, it will be sufficient to find out in which fluid do significant down-gradient wave fluxes of heat are observed to occur.

The zonal mode 2 scale of the ACW is selected in the present theory by surface drag on the low wavenumber side of the resonance of atmospheric Rossby waves stationary against the vertically averaged zonal wind. When the conditions are right, large growth rates can occur because the resonance is sharp (its width is controlled by the ratio of the advecting zonal wind time scale over the spin down time scale). Because of the sensitivity of this process to the vertical profile of the mean zonal wind, the optimal conditions for growth will be influenced by seasonal effects and the turbulent energy cascade of the transient eddies that bear upon the vertical structure of the mean zonal winds (of course the transient eddies effects are not limited to reshaping the vertical mean wind profile because the storm track modified by the large scale heating also acts to modify the lateral structure of the mean zonal flow). It is when the mean shear baroclinicity is large.ie. in winter that the wave number window for equivalent barotropic response is largest. Because of such seasonal variations, growth of the ACW is expected to be quite intermittent.

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7. Appendix

We show how the finite differencing in the vertical of the 2-level model equations that has been chosen ensures the conservation of total energy in the compressible case.

The internal energy equation is first written in flux form for each level i as:

$$\left\{\frac{\partial}{\partial t}\cdots + \nabla_{\mathrm{H}}\cdot \boldsymbol{u}_i \cdots\right\} \rho_i C_{\mathrm{p}} T_i + (-1)^i \frac{\rho_{\mathrm{I}} C_{\mathrm{p}} w_{\mathrm{I}}}{h_i} T_{\mathrm{I}} + \rho \tilde{w} T^i \frac{R}{H_0} = 0.$$

We write $\rho\tilde{w}T^i$ evaluated at level i as:

$$\rho\tilde{w}T^i = \rho_i T_{\mathrm{I}} \frac{w_{\mathrm{I}}}{2},$$

where $T_{\mathrm{I}} = \frac{1}{2}(T_1 + T_2)$ is a usual choice which guarantees conservation of variance for the incompressible part of the equations and $w_{\mathrm{I}}/2$ is the averaged of vertical velocities over the layer. Other choices are possible to conserve entropy Arakawa and Lamb, 1977 but for simplicity we restrict ourselves to conserve energy. Integrating over the domain (horizontal averaging is noted by brackets and vertical averaging by the sum $\sum$ over the layers) gives:

$$\frac{\partial}{\partial t} \sum \langle \rho_i C_{\mathrm{p}} T_i \rangle + \frac{R}{H_0} \left(\frac{\rho_1 h_1 + \rho_2 h_2}{2}\right) \langle w_{\mathrm{I}} T_{\mathrm{I}} \rangle = 0. \quad \text{(A1)}$$

The continuity equation is written as:

$$\rho_i \nabla_{\mathrm{H}} \cdot \boldsymbol{u}_i + (-1)^i \frac{\rho_{\mathrm{I}} w_{\mathrm{I}}}{h_i} = 0. \quad \text{(A2)}$$

After integrating over the domain the kinetic energy equation becomes:

$$\frac{\partial}{\partial t} \sum \left\langle \rho_i h_i \frac{|\boldsymbol{u}_i|^2}{2} \right\rangle = \rho_{\mathrm{I}} \langle w_{\mathrm{I}} (\phi_1 - \phi_2) \rangle. \quad \text{(A3)}$$

For total energy to be conserved the pressure work on the RHS of (A3) must be equal to the vertical flux term on the LHS of (A1). To be so ρ_{I} is defined as the height weighted density averaged over both levels:

$$\rho_{\mathrm{I}} = \frac{\rho_1 h_1 + \rho_2 h_2}{h_1 + h_2} \quad \text{(A4)}$$

and the form of the hydrostatic relation which guarantees equality of the two terms is found to be:

$$\phi_1 - \phi_2 = \frac{R}{H_0} T_1 \left(\frac{h_1 + h_2}{2} \right) \quad \text{(A5)}$$

The definition of (A4) and the form of (A5) ensure to deal with energy conserving equations. The decomposition, however, is non unique.

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