A Simple Mechanism for Blocking

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ABSTRACT

Numerical experiments show that blocking in a barotropic atmosphere can occur as a resonant enhancement of Rossby lee waves forced by two stationary sources of potential vorticity. In particular, if an upstream source of stationary forcing enhances the northerly flow over orography, then blocking occurs downstream of the mountain. In an analytical study, we show that, in the presence of friction, Rossby lee waves generate a rectified current downstream of the mountain, which does not vanish in the limit of zero friction.

The relevance of this study to observed generation of blocking in the Atlantic Ocean and immediately upstream of the Rockies is discussed.

1. Introduction

The discovery of baroclinic instability by Charney (1974) and Eady (1949) has dominated research in meteorology for the last 30 years. The large-scale atmospheric flow has been regarded as a baroclinically unstable circumpolar vortex, the evolution of which is sought in the presence of transient disturbances. While it has been recognized that the mean flow of the atmosphere is asymmetric as a result of fixed topographic and thermal asymmetries, most of the studies assumed the mean flow to be zonally homogeneous. General circulation models took into account the asymmetric nature of the boundary terms but they were not very successful in improving our basic understanding of the consequences arising from the asymmetry present. Yet such general circulation models did indicate that topography profoundly influences the behavior of the atmosphere (Manabe and Terpstra, 1974) as already had been known observationally for many years (Pettersen, 1956).

Among the atmospheric phenomena which are believed to be related to the interaction of traveling disturbances with asymmetric steady forcing is the blocking situation. Although the phenomenon had been known for many years (Rex, 1950a,b) the physical mechanisms underlying blocking remained obscure. Recently, Tung and Lindzen (1979) suggested that atmospheric blocking could be explained in terms of linear resonance of planetary-scale waves with respect to surface forcing such as continental elevation and land-sea differential heating. Egger (1978) proposed that blocking could be the manifestation of a barotropic nonlinear interaction among forced and slowly moving free waves. McWilliams (1980) has shown some similarity between vortex pair blockings and modon solutions of the barotropic equation. Charney and DeVore (1979) showed that an externally driven barotropic flow perturbed by topography and external wave-forcing possesses more than one nonlinear stable equilibrium state, one of which could be the blocking situation. While these latter two works suggest that nonlinearity is an essential ingredient of blocking, their use of zonal periodicity and truncated spectral representation cannot provide a clear-cut answer to the question whether such a prominent atmospheric phenomenon is local or global in character and what is the cause and effect relationship of the blocking situation.

With this motivation in mind we decided to perform numerical simulations under controlled conditions with the hope of obtaining sufficient clues that would lead to tractable analytical studies. We have chosen the barotropic vorticity equation as the starting point of our investigation, adopting the point of view (Charney and DeVore, 1979) that the interaction problem is dominated by energy transformations via wave-induced Reynolds stress fields while baroclinic processes are of secondary importance and are necessary only as far as triggering the traveling disturbances. Such an assumption seems particularly feasible in light of the pioneering works of Lorenz (1972) and Gill (1974) on the stability of barotropic
Rossby wave, which besides demonstrating that stability studies are not restricted to parallel flows, reinstated also the importance of barotropic processes for the dynamics of the atmosphere.

Our numerical experiments are performed for an open channel under prescribed upstream conditions. Removal of the constraint of zonal periodicity permits us to investigate more carefully the cause and effect relationship of the interaction problem. We find that when the system is repeatedly excited at some upstream location by localized disturbances which has a nonzero mean with respect to time, such as a succession of lows, the steady component of the response can dominate the time-dependent part of it, leading to an effectively time-independent field. The nonlinear interaction of this field with a preexisting steady asymmetric flow generated, for example, by localized topography, can lead to a new steady-state configuration.

The results show that depending on the nature of the excitation and on its phase relation with regard to the topography, two dominating steady flow configurations emerge: high-index flow with very little interaction between the two fields and a low index flow for which the interaction between the two fields is strongly nonlinear. This latter case corresponds to the blocking situation. It can be interpreted in a simple way using energy considerations which indicate that the nonlinear interaction is a consequence of a local finite-amplitude positive feedback mechanism. The new steady-state configuration can be maintained as long as the excitation lasts. When it ceases the system spins down to its original state. The excitation time scale and the spin-down time scale underly the persistent nature of the blocking configuration. For sufficiently small dissipation the system exhibits local oscillations immediately downstream of the topography. The oscillation can take place with or without time-dependent excitation. Some fundamental features of the basic steady state generated by isolated topography are investigated analytically in Section 2, where it is found that friction plays a singular role in the dynamics of barotropic Rossby waves. In Section 3 the design of the numerical experiments is described. The experimental integrations and the occurrence of blocking as a result of a nonlinear positive feedback mechanism are discussed in Section 4. A summary and further conclusions are presented in Section 5.

2. The basic state

The flow field generated by a westerly flow past an isolated topographic feature can be conceived as one prototype of a zonally asymmetric flow configuration. Since our goal is to understand the consequences of the interaction of such a field with disturbances excited upstream, this field should be adequately understood. Although a complete analytical investigation of the basic field is quite complicated, much relevant information can be obtained from the far-field solution which is readily obtainable. In this
section we summarize the three main results of such a study. The complete derivation is given in Appendix A.

We consider a steady, uniform westerly flow $U$ past an isolated topographic feature on a $\beta$ plane in a vertically bounded channel of width $L_y$ and depth $D$. When the channel is sufficiently narrow such that $L_y < \pi(U/\beta)^{1/2}$, where $\beta'$ is the gradient of the planetary vorticity, no stationary Rossby waves can exist and the response of the field to the topographic forcing is entirely local. This situation is depicted in Fig. 2.1a for an inviscid flow and in Fig. 2.1c for a slightly viscous flow. (Viscosity is assumed to affect the flow through Ekman suction only.) In the inviscid case the flow response consists of a single high-pressure cell centered over the topography. In the viscous case the high-pressure cell is shifted upstream of the topography and a weak trough appears in the lee. We emphasize that the viscous solution depicted in Fig. 2.1c is valid only for weak topographic forcing. In contrast, the inviscid solution depicted in Fig. 2.1a is a finite-amplitude solution provided no closed circulations are present in the flow field.

The character of the flow field changes drastically when the channel is sufficiently wide, i.e., $L_y > \pi(U/\beta')^{1/2}$ such that stationary Rossby waves can be excited. If $\pi < L_y(\beta'/U)^{1/2} < (n + 1)\pi$, where $n$ is a positive integer, $n$ stationary Rossby waves can be excited. The corresponding wavenumbers are given by $[(\beta'/U) - (n\pi/L_y)]^{1/2}$. Fig. 2.1b depicts the inviscid steady-state solution when only one forced Rossby wave is excited. We observe that the topography does not affect the flow field on its wind side but that a forced Rossby wave is excited in the lee side. Again, this solution is not restricted to weak topographic forcing as long as no closed circulations appear in the flow field.

Friction damps out the Rossby lee wave and the effect of this is to generate rectified currents. These currents do not vanish in the limit of zero friction implying that the strictly inviscid solution shown in Fig. 2.1b is not realizable. The effect of the currents is to cause a tilt of the Rossby lee wave as exemplified in Figs. 2.1d–2.1f. The direction of the rectified current, and therefore the tilt of the waves, depends on whether $L_y(\beta'/U)^{1/2}$ is larger, smaller or equal to 2 (Figs. 2.1d, 2.1e and 2.1f, respectively). The case in Fig. 2.1d (wider channel) is particularly interesting since the tilt of the rectified current is such that it resembles an omega blocking pattern. As stated before, the complete analysis leading to these results is included in Appendix A.

3. Design of the numerical experiments

The experiments that guided this study were performed with physical scales relevant to midlatitude atmospheric flow. We integrated the quasi-geostrophic potential vorticity equation, which in dimensional form is

$$[\partial/\partial t + (U + v') \cdot \nabla] \eta' = -\nabla^2 \psi'/\tau,$$  \hspace{1cm} (3.1)

where $U$ is the basic uniform flow (10 m s$^{-1}$ in most experiments), $v' = k \times \nabla \psi'$ is the perturbation flow, $\nabla^2 \psi' = k \cdot \nabla \times v'$ is the perturbation relative vorticity and $\tau$ the frictional time scale. The potential vorticity is defined by

$$\eta' = \beta y + \nabla^2 \psi + f_0 h/D,$$ \hspace{1cm} (3.2)
a summation of the planetary and relative vorticity, and of the orographic component due to vertical compression and stretching of vortex tubes.

We performed the integration on a very long open channel (Fig. 3.1), with a 32 000 km zonal extent. This was done in order to study the local effect of isolated orography avoiding the conceptual complication of lee wave reentering a periodic channel. A single narrow mountain defined by

$$h/D = \alpha \sin(\pi y/L_y) \exp[-(x/\Delta x)^2]$$

was located at the center of the channel, where $D$ is the mean depth of the atmosphere, $h$ the height of the mountain and $\alpha$ its maximum relative amplitude. The flow was driven by a constant unperturbed zonal flow $U$ upstream.

The numerical scheme, based on enstrophy conserving space differences combined with Lorenz' $N$-cycle scheme, and the boundary conditions are described in the Appendix B.

After the steady-state solution is established, no energy propagates upstream from the orography, because waves with zero phase speed have positive group velocity. Nevertheless, during the transient stage some energy propagates upstream and, because of the boundary conditions of unperturbed zonal flow at the entry of the channel, they cannot leave the domain and become reflected. The reflected waves interact again with the mountain, and set up further transient waves that move upstream and become reflected. This process acts resonantly, with a slow amplification of the waves spuriously reflected at the entrance of the channel, and makes the simulation meaningless (Fig. 3.2). It is not unlike the situation in primitive equation atmospheric models which use $\omega = 0$ as a top boundary condition. This artificial condition, like the one used in our experiment, is well posed but leads to spurious reflections.

We solved the problem with a sponge layer 10 000 km long at the upstream end of the channel that absorbs the energy of incident waves. Following Cardelino (1978), we used a friction coefficient $\tau$ that varied linearly from (2 days)$^{-1}$ at the entrance to the value (20 days)$^{-1}$ used in the interior (Fig. 3.1).

The sponge layer eliminated completely the problem of spurious reflection (Fig. 3.3), without affecting the steady solution which has no waves traveling upstream.

This work was started with the purpose of studying the interaction between stationary Rossby waves excited by the presence of orography and transient Rossby waves. The latter were simulated introducing perturbations in the form of "pulses" or "lows" by the following procedure: Every four days we increased the instantaneous value of the potential vorticity field over a region of size $3\Delta x$ by $3\Delta y$ (a square of 1500 km side) by an amount equal to 30% of the mean planetary vorticity $f_0$. These pulses or lows were introduced at a distance upstream from the mountain that varied between 2000 and 8000 km in the different experiments. In some experiments we generated highs by adding negative vorticity, or alternated highs and lows every two days. Finally, we also performed experiments in which the source of pulses was replaced by steady forcing in the form of a Gaussian mountain defined by $h/D = 0.3 \exp[-(x^2 + y^2)/(500 \text{ km})^2]$.

4. Experimental results

a. Generation of perturbation kinetic energy

In the interpretation of the numerical results, it is very helpful to consider the equation for the perturbation kinetic energy (KE'). If we multiply equation (3.1) by $-\psi'$, integrate over the spatial domain, and make use of Gauss' theorem, we obtain

$$\frac{\partial}{\partial t} \left[ \int_{-L_y/2}^{L_y} \int_0^{L_x/2} \frac{1}{2} \left| \nabla \psi' \right|^2 \, dx \, dy \right] = -Uf_0 \int_{h\pi_0}^{L_y} (h/D)v' \, dx \, dy$$

$$- \int_{-L_x/2}^{L_x/2} \int_0^{L_y} \nabla \psi' \cdot \nabla (\psi'/\tau) \, dx \, dy$$

$$- \frac{1}{2} U \int_0^{L_y} \left[ (\partial \psi'/\partial x)^2 \right]_{-L_x/2}^{L_x/2} \, dy,$$

Fig. 3.3. As in Fig. 3.2 except with a sponge layer.
which indicates that the time variation of $KE'$ is given by the sum of three terms. The first one is the generation of $KE'$ by the interaction between the mean flow $U$, the perturbation field $\psi'$ and the orography, the second one is the dissipation of $KE'$, and the third is the advection of $KE'$ through the open ends of the channel. Of these terms, it is noteworthy the simple expression of the generation term which indicates that energy extraction from the mean flow occurs only over topography, and requires northerly perturbation flow ($v' < 0$) over the region with positive orography ($h/D > 0$). In other words, the pressure gradient force $-\partial \psi'/\partial x$ over the mountain has to be in the same direction as the mean flow. This implies that if there is no generation of $KE'$ then the pressure field must be in phase with the orography. Under steady-state conditions, this can only occur if there is no friction and if $\beta = 0$, so that no $KE'$ is dissipated or advected away from the region of generation (cf. Fig. 2.1a). If, on the other hand, $KE'$ is being generated, the pressure wave must be shifted upstream, with high pressure on the wind side and low pressure on the lee side of the mountain. (Figs. 2.1b and 2.1f.). This analysis of the generation terms is not limited to forcing by orography. The effect of any type of stationary forcing (orography, heating, even locally enhanced baroclinic instability) can be represented as a source of vorticity in the potential vorticity equation (3.1) of the form $\nabla \cdot \nabla \psi_0 / T$, where $T$ is a time scale related to the forcing mechanism. With this expression for stationary forcing, the $KE'$ generation term in equation (4.1) becomes

$$G = T^{-1} \int \int \nabla \psi_0 \cdot \nabla \psi' \, dx \, dy,$$  

(4.2)

where the integral is performed in the region where $\psi_0$ is not constant. Again, we see that the sign of the generation of $KE'$ depends on the phase relationship between the perturbation flow and the streamlines associated with the vorticity source.

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**b. Experiments with a single mountain**

The approximately steady-state solutions obtained with a single topographic feature in the center of the channel agree well with the analysis presented in Section 2.

Fig. 4.1 shows the solution obtained with our standard set of parameters: $U = 10$ m s$^{-1}$, $L_y = 5000$ km, $(h/D)_{\text{max}} = 0.2$ (corresponding to a mountain of maximum height of 2000 m), and a frictional time scale $\tau = 20$ days. The solution is attained after the initial transient is washed away from the domain of integration ($\sim 20$ days). There is a very small-amplitude oscillation, as can be observed in the generation of perturbation kinetic energy (Fig. 4.8, labeled no pulses). The oscillation, which is due to the rather small value of the friction, disappears if we use $\tau = 10$ days. This case corresponds to Fig. 2.1d. The backward tilt of the Rossby lee wave due to the frictionally induced rectified current discussed in Section 2 can be observed in the far field downstream of the mountain. It should be noted that the phase of the pressure field is such that there is maximum generation of perturbation kinetic energy, i.e., $-\partial \psi'/\partial x$ is maximum at $x = 0$. This phase relationship was present in all our experiments with a single source of lee waves.

Fig. 4.2 exhibits the solution when the amplitude of the mountain is increased to the equivalent of 3000 m. It is clear that such a single large-amplitude localized source of vorticity can produce "blocking" Rossby lee waves. In Fig. 4.3 the meridional width of the channel has been reduced to 2500 km, eliminating in this way the existence of Rossby lee waves (as in Fig. 2.1c). Even in this case we see that the maximum zonal pressure gradient occurs at the center of the channel. In Fig. 4.4 the meridional width of the channel has been increased to 8000 km, allowing for meridional wavenumbers $n = 1, 2$ and 3. Since the mountain forces only wavenumber 1, it is this wavenumber that is apparent in the solution.

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**Fig. 4.2.** As in Fig. 4.1 except with $(h/D)_{\text{max}} = 0.3.$
c. Experiments with pulses generated upstream of the mountain

We performed a large number of experiments in which "pulses" of positive (negative) vorticity simulating "lows" ("highs") were introduced upstream of the mountain, as described in Section 3. In different experiments we varied the frequency at which we sent the pulses (once every 2–4 days), and the distance upstream from the mountain (2000–10 000 km). In the standard case \( L_v = 5000 \) km we observed that the numerical integrations reached a new, almost constant circulation 10–20 days after the pulses were started.

The result corresponding to the standard case with positive pulses (lows) sent from \( x = -6000 \) km every three days is shown in Fig. 4.5. This result is typical of most of the experiments; there are small-amplitude waves downstream of the region of generation of the lows (indicated by a shaded square) but almost no noticeable change in the mountain lee waves (cf. Figs. 4.1 and 4.5).

However, when the pulses were sent from certain distances the solution changed dramatically. Fig. 4.6 corresponds to the same experiment as Fig. 4.5 except that the pulses were launched from \( x = -4000 \) km. The waves upstream of the mountain are very weak. In the lee of the mountain, on the other hand, there is an omega blocking type of circulation, with a pronounced low and high and a splitting of the jet.

We observed that the generation or lack of generation of a blocking high downstream of the mountain was rather insensitive to the frequency of the pulses. When we sent alternating lows and highs every two days, the effect of the pulses disappeared and we recovered the unperturbed solution. In order to explain these results, we performed integrations in which the pulses were generated in the absence of orography. The result was a small-amplitude time-averaged flow of lee waves downstream of the region of generation (Fig. 4.7) superimposed with a much smaller amplitude time-varying solution. It is the time-averaged solution that can interact resonantly with the orography downstream so that our results are independent of the frequency of generation. This was confirmed by replacing the pulses with a small upstream Gaussian mountain as indicated in Section 3, and observing similar results.

We verified that two conditions must be met in order to develop blocking solutions.

First, the position of the upstream forcing region must be such that its lee waves produce maximum generation of \( KE' \) when they interact with the mountain downstream. In other words, the lee waves corresponding to the upstream region of forcing must have \(-\partial \psi'/\partial x\) large and positive as they cross the downstream mountain. The resonant interaction occurs because the source of positive vorticity produces a weak ridge centered at \(-2800 \) km downstream of it, and a weak trough at \(-3600 \) km, the wavelength of the stationary Rossby wave for the given parameters. When the distance between two sources is \( 4000 \) km, or about \( \frac{3}{4} \) of the wavelength of the stationary Rossby wave, the mountain is between the ridge and the trough induced by the upstream source of vorticity. The combined forcing is very strong at the wavelength of stationary Rossby wave, and a large blocking response is obtained (Fig. 4.6).

If the mountain is located elsewhere, as in the case depicted in Fig. 4.5, in which the mountain is slightly ahead of the source's trough, then no resonant forcing is produced, and the response is dominated by the mountain forcing.

Second, we found that the amplitude of all forced fields must be sufficiently large. When we reduced the amplitude of either forcing by a factor of 2, no blocking was observed. This indicates that the blocking phenomena are not due to a small-amplitude instability mechanism.

The relationship of the distance between the re-
gions of forcing and generation of $KE'$ can be observed in Fig. 4.8, corresponding to the three cases of no pulses, no blocking ($x = -6000$ km), and blocking ($x = -4000$ km). Note the rapid enhancement of generation of $KE'$ after the signal of the pulses from $x = -4000$ km first reaches the orography. As indicated before, the generation of $KE'$ is insensitive to the frequency of the pulses. The initial overshooting in generation is typical of all the experiments.

Our experiments suggest that in a barotropic fluid with isolated orographic forcing, blocking is a phenomenon which is locally generated through a nonlinear positive feedback mechanism. When the flow over the mountain has a northerly component, there is generation of eddy kinetic energy which is realized downstream of the mountain as Rossby lee waves. If the northerly flow is enhanced, for example, by upstream forcing occurring at a favorable position, then the intensity of the first low downstream of the mountain increases. This, in turn, strengthens the northerly flow over the orography, and therefore enhances the generation of eddy kinetic energy. The growth of $KE'$ continues until the generation of $KE'$ is balanced by dissipation. If the frictional dissipation is small, there is overshooting in the generation, and we observe vacillation. The process that gives rise to vacillation also is local in nature: the generation of $KE'$ deepens the low in the lee of the mountain, increasing therefore the generation of $KE'$. The perturbations are overamplified, until advection of $KE'$ out of the domain reduces the intensity of the low and the generation of $KE'$. When the advection of $KE'$ weakens, the process can start again.

It is important to note that blocking due to enhancement of generation of $KE'$ can take place under any kind of forcing and not just orography. The only condition for enhanced generation of $KE'$ is that the flow over the region where there is stationary forcing must be in phase with the streamlines implied by the forcing [Eq. (4.2)].

We also performed experiments in which we stopped the upstream forcing (pulses) after the blocking circulation had been established. We observed that once the excitation ceased, blocking persisted for several days. The flow pattern returned to its high-index type of unperturbed circulation over a period of the order of 20 days, the frictional time scale. This observation may help explain the persistent character of blocking.

d. Weakly nonlinear experiments

The realistic blocking development obtained when the upstream perturbation enhanced the generation of $KE'$ over the orography (Fig. 4.6), is a resonance phenomenon. In order to study the importance of nonlinear interactions, we repeated the experiments of the previous subsection integrating the equation in which the only nonlinear term retained is the interaction of the perturbation flow and the orography. Figs. 4.9a, 4.9b and 4.9c correspond to the integrations without pulses, with pulses generated at $x = -6000$ km and at $x = -4000$ km, respectively. In the latter case there is a strong resonant enhancement. However, the pattern in Fig. 4.9c does not resemble blocking. It should also be noted that the effect of a rectified current downstream, discussed in Section 2 and clearly observable in Figs. 4.4 to 4.6 is absent in Fig. 4.9. The nonlinear self-interaction is necessary to produce the omega-type of blocking of Fig. 4.6.

e. Blocking in a wider channel

Similar experiments were performed on a wider channel, with $L_y = 8000$ km, allowing Rossby lee waves with meridional wavenumber 1, 2 and 3 to be excited. Fig. 4.10 is the solution when pulses of negative vorticity (highs) are sent from $x = -4000$ km, and is quite similar to the solution without pulses (Fig. 4.4). Negligible changes also were observed.
when positive pulses (lows) were sent from other distances upstream of the mountain. When lows were sent from \( x = -4000 \) km, however, we observed a very pronounced blocking pattern (Fig. 4.11).

In this case we also observed two other important changes. First, the solution remained strongly time dependent, with the blocking high downstream of the mountain showing a vacillation amplifying with time. Second, a strong excitation of meridional wavenumber 2 occurred downstream of the pulses but immediately upstream of the mountain. It appears as a strong high developing at northern latitudes with lower pressure at midlatitudes. The fact that a similar “blocking high” upstream of the mountain appears also in weakly nonlinear integrations indicates that it is due to nonlinear interaction between the lee waves of the perturbation upstream and the orography, rather than to self interaction of the perturbation flow.

5. Summary and discussion

We have shown that blocking in a barotropic atmosphere can occur as a result of resonant enhancement of Rossby lee waves forced by two stationary sources of potential vorticity.

The structure of the stationary response to one isolated mountain in an infinite \( \beta \)-plane channel has been discussed in an analytical study of the far-field solution. We showed that friction plays a singular role in the dynamics of Rossby lee waves. It generates a rectified current which does not vanish in the limit of zero friction. In a wide channel, the rectified current introduces a tilt in the Rossby lee waves that resembles blocking.

An analysis of the energy equation indicates that the generation of Rossby lee waves depends on the phase of the flow over a region of forcing. In the case of orographic forcing, for example, generation of eddy kinetic energy is proportional to the intensity of northerly flow over the mountain. We performed numerical integrations of barotropic flow in an open \( \beta \)-plane channel. The most important result of the numerical experiments is the following: If an upstream source of stationary forcing enhances the northerly flow over a mountain, then blocking occurs downstream of the mountain. Experiments performed with a quasilinear version of the barotropic potential vorticity equation also show resonant enhancement of the Rossby lee waves. However, the full nonlinear interactions are necessary in order to generate realistic omega-type of blocking patterns.

When the experiments are performed in a channel wide enough to allow Rossby lee waves with more than one meridional wavenumber, a new result is obtained. In this case, an upstream source of stationary forcing with a favorable phase generates blocking not only downstream, but also immediately upstream of the mountain. Weakly nonlinear experiments demonstrate that the upstream blocking high, which occurs at high latitudes only, is a result of the interaction between the response to the upstream forcing and the mountain.

We believe that our results may be relevant to ob-
served atmospheric blocking. Sources of stationary or quasi-stationary forcing are not just orography, but sea surface temperature anomalies and even regions of enhanced baroclinicity, where repeated cyclogenesis may occur. Blocking can take place through enhancement of Rossby lee waves if two such regions are in a favorable phase. This local resonance may occur only during certain periods because of the dependence of the wavelength of stationary Rossby waves on the intensity of the mean zonal flow.

According to Rex (1950a,b), blocking highs occur in the Atlantic and Pacific oceans east of the great semipermanent oceanic cyclonic centers (the Icelandic and Aleutian lows, respectively). Rex (1950b) analyzed the seasonal trend in 82 Atlantic cases of blocking. He found (cf. his Fig. 2) that the Atlantic blocking center always occurs at a longitude 40°–60° east of the Atlantic low. The semipermanent cyclones are a result of stationary forcing related to orography and/or land-sea-ice contrast. The fact that many cyclone tracks end in the semipermanent lows also may be due to boundary forcing. If our barotropic experiments are relevant, we may infer that blocking can occur when the time-averaged flow coming from the North American continent is in phase with the Icelandic low and, consequently, enhances the generation of perturbation kinetic energy over the region of forcing.

Observational studies indicate that blocking frequently appears as a train of damped lee waves, similar to our results (Rex, 1950a, Figs. 1 and 2). Dole (1981) has observed that composite time-averaged pressure anomaly fields during blocking also show a similar structure. The monthly averaged 700 mb maps during the last 10 winter seasons suggest that persistent blocking in the Atlantic is associated with anomalously strong northerly flow over the Rockies. This is particularly true in the persistent anomalous circulation observed in December 1976 (Taubensee, 1977) and January 1977 (Wagner, 1977). Indeed, enhanced orographic gen-

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eration of perturbation kinetic energy seems to have been associated with a favorable distribution of sea surface temperatures in the Pacific Ocean. Following a suggestion of Namias (1978), Shukla and Bangaru (1979) have recently performed a sensitivity experiment with a GLAS general circulation model in which they introduced observed sea surface temperature anomalies in the Pacific multiplied by a factor of 9/5 (Fig. 5.1). The sea surface temperature anomalies enhanced the generation term over the Rockies (Fig. 5.2) and in good agreement with our simple barotropic model, produced a maximum response downstream of the regions of orographic and thermal forcing.

The high-latitude blocking high upstream of the mountain observed in our experiment with a wide channel also appears to have an atmospheric counterpart. It resembles the blocking highs that sometimes develop west of the Canadian Rockies, accompanied by low pressure in the coast of California.\(^2\) According to our experiments this may be due to an interaction of the atmospheric response to anomalous Pacific sea surface temperatures with the Rockies.

The fact that the generation of perturbation kinetic energy by orography is proportional to the northerly flow over the mountains suggests the existence of a positive feedback mechanism. Strong generation of perturbation kinetic energy will intensify the low in the lee of the mountain, thereby increasing the zonal pressure gradient over the orography. This, in turn, increases the northerly flow and the generation of eddy kinetic energy. The positive feedback mechanism allows a simple interpretation of the multiple flow equilibria found by Charney and DeVore (1979) for nonlinear flow over topography in a periodic $\beta$-plane channel, and by Källén (1980) in spherical geometry. Under initial conditions corresponding to high-index flow, $v$ is small over orography, the generation of eddy kinetic energy is inefficient, and the flow remains in the high index "attraction basin". If the initial flow has a large-amplitude perturbation (low index) with a favorable phase, there will be strong generation, and the flow will remain low index. If the phase is

\(^2\) J. Wallace, personal communication.
unfavorable, eddy kinetic energy will decay and the flow will evolve toward high index.

Hartmann and Ghan (1980) have recently studied the local balance in the vorticity and temperature equations for blocking highs and transient ridges in the Atlantic and Pacific oceans. They found that with the obvious exception of smaller time derivatives in the case of blocking highs (which they traced to nonlinear terms), the local characteristics of the persistent and transient phenomena were remarkably similar. Our experiment suggests that the reason for this result is that blocking is not locally generated, but its cause must be found upstream. We should note, however, that a "modon" type of solution (McWilliams, 1980) also may be consistent with Hartmann and Ghan's results.

We have emphasized the importance of Rossby lee waves in blocking. In a recent study of the atmospheric circulation during the FGGE period 5–21 January 1979, Kalnay-Rivas et al. (1980) have found the presence of very large-amplitude, short wavelength, Rossby lee waves in the Southern Hemisphere. They appear downstream of the Andes suggesting that in the Southern Hemisphere, where orographic and land-sea contrasts occur in very localized regions, Rossby lee waves may be an important component of the general circulation. Observed blocking in the Atlantic region east of South America may be related to the presence of these waves.

We plan to continue this study with a two-level baroclinic model, and to introduce spherical geometry. We also will study the relationship between the flow over the Rockies and other regions of forcing and the persistence of ridges in the Atlantic. It is hoped that observational studies guided by simple theoretical studies such as ours may provide us with a better understanding of atmospheric blocking.

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APPENDIX A

Basic-State Solution

We consider a steady uniform westerly flow $U$ past an isolated topographic feature on a beta plane in a vertically bound channel of width $L$ and depth $D$. The nondimensional quasi-geostrophic vorticity equation governing the perturbation streamfunction is

$$\nabla^2 \psi + \beta \phi_x + r \nabla^2 \psi + (h_0/Ro)F_x = -J(\psi, \nabla^2 \psi + (h_0/Ro)F), \quad (A1)$$

where the nondimensional parameters are defined as

$$\begin{align*}
Ro &= U/(f_0L) < O(1) \\
h_0 &= h/D \ll O(Ro) \\
\beta &= \beta' L^2/U \ll O(1) \\
r &= L(vf_0)^{1/2}(DU) \ll O(1)
\end{align*}$$

$h$ is the height of the topography and $F(x, y)$ its functional shape, considered here to be positive; $f_0$ is the value of the Coriolis parameter at the latitude about which the beta plane approximation is made. $\beta'$ is the dimensional gradient of the earth's vorticity at that latitude. The horizontal coordinates are scaled with $L$ where $x$ points eastward. The non-dimensional streamwise extent of the topography is $a$ and we consider the case $a \ll O(1)$. $U$ is the velocity scale and the dimensional streamfunction for the unperturbed current is given by $-Uy$. Eq. (A1) is supplemented by the boundary conditions

$$\begin{align*}
\psi &= 0 \text{ on } y = 0, 1 \\
\psi &= 0 \text{ as } x \to -\infty \quad \text{(radiation condition)} \\
\psi \text{ is bounded for } x \to \infty
\end{align*} \quad (A2)$$

In the next section we determine the linear and nonlinear characteristics of the far-field solution.

1. Inviscid solution

In the absence of closed streamlines and without linearization it can be shown that the perturbation streamfunction is governed by

$$\nabla^2 \psi + \beta \phi_x + (h_0/Ro)F = 0 \quad (A3)$$

subject to the boundary conditions (A2). Away from the topography (A3) becomes

$$\nabla^2 \psi + \beta \phi_x \sim 0 \quad \text{for } |x| \gg a. \quad (A4)$$

The meridional structure of the disturbance is determined by the projection of the meridional structure of the topography on the meridional eigenfunctions, viz., \sin(n\pi y), \; n = 1, 2, \ldots. \; \text{Let } k \text{ be a zonal wavenumber satisfying the dispersion relation}

$$k^2 = \beta - n^2 \pi^2, \quad (A5)$$

then for $\beta < \pi^2$ (A5) yields only evanescent

$^3$ The analysis presented here is more easily performed using nondimensional equations.
standing waves. It is not difficult to show that for a
topography with $\sin \pi y$ meridional structure the flow
response consists of a single high-pressure cell
centered over the topography. The asymptotic far-
field response is given by

$$\psi \sim \frac{h_0}{Ro} C \exp(-(\pi^2 - \beta)^{1/2}|x|) \sin \pi y,$$

$$|x| \gg a,$$  \hspace{1cm} (A6)

where $C$ is a positive constant.

$N$ Rossby lee waves are excitable provided
$N^2 \pi^2 < \beta < (N + 1)^2 \pi^2$. The actual number of the
lee waves excited depends on the meridional struc-
ture of the topography. The asymptotic far field of
the first Rossby lee wave is given by

$$\psi \sim \frac{h_0}{Ro} D \cos((\beta - \pi^2)^{1/2}x + \alpha) \sin \pi y$$

$$x \gg a$$  \hspace{1cm} (A7)

$$\psi \sim 0, \hspace{0.2cm} x \ll -a$$

The phase $\alpha$ and the sign of $D$ should be consis-
tent with the existence of a lee trough immediately
downstream of the topography. As $N$ increases the
additional lee waves have longer zonal wavelengths
but they possess more wiggles in the meridional
direction.

2. Viscous solution

Unlike the inviscid case, the vorticity equation
(A1) cannot be reduced to a linear equation which
implies that finite-amplitude closed-form solutions
cannot generally be obtained. However, much infor-
mation can be extracted from weakly nonlinear
analysis valid for $(h_0/\eta) = \epsilon \ll 1$. The linearized
vorticity equation becomes

$$\nabla^2 \psi_x + \beta \psi_x + r \nabla^2 \psi + \epsilon F_x = 0,$$  \hspace{1cm} (A8)

with the far-field behavior governed by

$$\nabla^2 \psi_x + \beta \psi_x + r \nabla^2 \psi \sim 0, \hspace{0.2cm} |x| \gg a.$$  \hspace{1cm} (A9)

Note that (A9) also is a proper asymptotic limit for
$\epsilon = O(1)$.

In order to simplify our discussion as much as
possible we assume that the only meridional struc-
ture of the topography is $\sin \pi y$ and consider sepa-
rately the two cases $\beta < \pi^2$ and $\beta > \pi^2$. Eq. (A9)
possesses three independent solutions of the form
$e^{ikx} \sin \pi y$, where $k$ satisfies the dispersion relation

$$k^3 - ikr^2 + k(\pi^2 - \beta) - ir\pi^2 = 0.$$  \hspace{1cm} (A10)

A geophysically relevant case is $r < O(1)$ for which
we obtain that the three roots of (A10) are approxi-
mated by

$$k \sim \pm i(\pi^2 - \beta)^{1/2} - \frac{i\beta r}{[2(\pi^2 - \beta)]}$$
$$k \sim i\pi^2 r(\pi^2 - \beta)$$

for $\beta < \pi^2$ and $r \to 0$ while the corresponding
approximations for $\beta > \pi^2$ and $r \to 0$ are

$$k \sim \pm(\beta - \pi^2)^{1/2} + i\beta r/[2(\beta - \pi^2)]$$
$$k \sim -i\pi^2 r(\beta - \pi^2)$$

(A11)

(A12)

When $\beta < \pi^2$ the far-field solution is given by

$$\psi \sim C_1 \epsilon \exp((\pi^2 - \beta)^{1/2}x$$
$$+ \frac{1}{2} \beta \pi^2 x/(\pi^2 - \beta)) \sin \pi y,$$

$$x \ll -a$$  \hspace{1cm} (A13)

$$\psi \sim C_2 \epsilon \exp(-(\pi^2 - \beta)^{1/2}x$$
$$+ \frac{1}{2} \beta \pi^2 x/(\pi^2 - \beta)) \sin \pi y$$

$$+ C_3 \epsilon \exp(-\pi^2 r x/(\pi^2 - \beta)) \sin \pi y,$$

$$x \gg a$$

In the limit of vanishing viscosity the inviscid solu-
tion should be recovered which implies that

$$C_1 = O(1), \hspace{0.2cm} C_2 = O(1), \hspace{0.2cm} C_3 = o(1)$$

as

$$r \to 0.$$  \hspace{1cm} (A14)\footnote{O(1) represents "order of magnitude smaller than."}

In accord with the inviscid limit both $C_1$ and $C_2$
should be positive. $C_3$ is negative as the follow-
ing considerations show. The fluid column which
ascends the topography is compressed acquiring nega-
tive relative vorticity which is somewhat reduced by
Ekman spin-up. The spin-up of negative vorticity con-
tinues when the fluid column descends the topo-
graphy, but now vortex stretching, reduced some-
what by Ekman spin-down, is also taking place. At the
bottom of the topography the fluid column is left with
positive vorticity which is spun-down downstream.
Since the downstream vorticity is positive and since
at distances $O(1/r)$ the second term of (A13) for
$x \gg a$ is dominant, although it is negligible com-
pared to the first term for $x \ll O(1)$, it follows that
$C_3 < 0$.

When $\beta > \pi^2$ the far-field solution is given by

$$\psi \sim B_1 \epsilon \exp(\pi^2 r x/(\beta - \pi^2)) \sin \pi y,$$

$$x \ll -a$$  \hspace{1cm} (A15)

$$\psi \sim B_2 \epsilon \exp(-\frac{1}{2} \beta \pi^2 r x/(\beta - \pi^2))$$
$$\times \cos((\beta - \pi^2)^{1/2} x + \alpha), \hspace{0.2cm} x \gg a$$

($\beta/[2(\beta - \pi^2)]$ is the reciprocal of the group velocity
of the Rossby lee wave evaluated in the limit of
$r \to 0$.) In the limit of vanishing viscosity the in-
viscid solution should be recovered which implies

\hspace{1cm}
that
\[ B_1 = O(1), \quad B_2 = O(1) \quad \text{as} \quad r \to 0. \quad (A16) \]
The effect of friction, besides damping the Rossby lee wave, is to induce a region of negative vorticity and consequently high pressure on the windward side of the topography. This is a consequence of vorticity continuity required by the viscous model and the existence of a region of negative vorticity over the topography. It follows that \( B_1 \) is positive. Again, the phase \( \alpha \) and the sign of \( B_2 \) should be consistent with the existence of a lee trough immediately downstream of the topography.

Since the linear viscous solution does not satisfy the full nonlinear equation an estimate of the neglected terms is warranted. Away from the topography the full vorticity balance is given by
\[ \nabla^2 \psi_x + \beta \psi_x + r \nabla^2 \psi = -J(\psi, \nabla \psi), \quad (A17) \]
where it is assumed that \( \psi = O(\varepsilon) \) and \( \varepsilon \ll 1 \). Expanding the perturbation streamfunction \( \psi \) in powers of \( \varepsilon \), viz.,
\[ \psi \sim \epsilon \psi' + \epsilon^2 \psi^2 = \cdots \quad (A18) \]
we find that for \( \beta > \pi^2 \) the \( O(\varepsilon) \) far-field solution is given by \( (A15) \). The nonlinear correction is governed by
\[ \nabla^2 \psi_x^2 + \beta \psi_x^2 + r \nabla^2 \psi = \frac{\beta B_2}{2(2\pi^2 - \beta)} \frac{\beta - \pi^2}{2\pi^2} \left( s \nabla^2 \psi \right), \quad (A19) \]
where we are considering only the downstream field. There is no \( O(r) \) nonlinear correction for the upstream field and similarly for the case at \( \beta < \pi^2 \). The \( O(\varepsilon^2) \) correction can now be determined and we obtain
\[ \psi^2 = \frac{\beta B_2}{2(2\pi^2 - \beta)} \frac{\beta - \pi^2}{2\pi^2 + \beta} \left( 1 + O(r) \right) \left( \psi \right)^2 \sin(2\pi y) \exp[-\beta r x/(\beta - \pi^2)], \quad x \gg a. \quad (A20) \]
It follows from \( (A20) \) that the \( O(\varepsilon^2) \) correction to the perturbation streamfunction does not vanish in the limit of \( r \to 0 \) although the Jacobian of the \( O(\varepsilon) \) field [the right-hand side of \( (A19) \)] is proportional to \( r \). Note that in the strictly inviscid case, viz., \( r = 0 \) the Jacobian of the \( O(\varepsilon) \) field vanishes identically indicating that the linear solution satisfies exactly the nonlinear equation. We conclude that there exists nonuniformity of the limits \( \varepsilon \to 0 \) and \( r \to 0 \). The effect of friction is singular and it is manifested in those parts of the flow field which support spatially attenuating Rossby waves. The cause for the non-uniformity can be traced to the dispersion relation. The inviscid dispersion relation is degenerate since it is satisfied by any steady zonal flow. This degeneracy is removed with the inclusion of friction no matter how small. Thus friction plays a singular role in the dynamics although it does not formally raise the order of the equation.

Eq. \( (A20) \) leads to the interesting result that spatial attenuation of Rossby waves generates currents—a conclusion which is valid also for traveling Rossby waves. This mechanism should be contrasted with the zonal flow correction generated by baroclinic instability (Pedlosky, 1970), or by nonlinear interaction of resonating Rossby waves (Lösch, 1977). It is evident from \( (A20) \) that \( \beta < 2\pi^2 \) or \( \beta > 2\pi^2 \) yield qualitatively different nonlinear corrections. In both cases, however, the nonlinear correction is \( O(\varepsilon^2) \) uniformly. The situation changes when \( \beta = 2\pi^2 \) since the denominator of \( (A20) \) vanishes suggesting a large nonlinear response. We propose to show that this conclusion is false. For \( \beta = 2\pi^2 \) the flow field exhibits frictional induced nonlinear resonance phenomenon. The nonlinear correction is proportional to \( \epsilon r \) and, consequently, vanishes when \( r \to 0 \).

Thus for \( \beta = 2\pi^2 \) the uniformity of the limits \( \varepsilon \to 0 \) and \( r \to 0 \) is recovered indicating that the linear solution is an exact solution of the nonlinear equation. It can be verified by direct substitution that for \( \beta = 2\pi^2 \) the right-hand side of \( (A19) \) satisfies the linear operator on the left-hand side with \( O(r^2) \) relative error. Thus for \( \beta = 2\pi^2 \) and \( x \ll 1 \exp(-2r x) \times \sin(2\pi y) \exp(-2r x) \) which, indeed, vanishes in the limit \( r \to 0 \).

The secular behavior of \( x \) is not dangerous here because of the exponential decay. Nevertheless, a more careful analysis is desirable if the possible modulation of the \( O(\varepsilon) \) solution is to be accounted for. The analysis is modified by allowing for long-scale modulation in the perturbation streamfunction which is written now as
\[ \psi \sim \epsilon \psi'(x, y) + \epsilon^2 \psi^2(x, y) + \cdots; \quad \chi = \epsilon x. \quad (A21) \]
The linear problem is left unchanged but the next correction is governed by
\[ \nabla^2 \psi_x^2 + \beta \psi_x^2 + r \nabla^2 \psi = \frac{\beta B_2}{2(2\pi^2 - \beta)} \frac{\beta - \pi^2}{2\pi^2 + \beta} \left( \psi \right)^2 \sin(2\pi y) \exp[-\beta r x/(\beta - \pi^2)], \quad (A22) \]
The \( O(\varepsilon) \) solution can be written as
\[ \psi' = \frac{1}{2}(A(\chi)e^{ikx} + A^*(\chi)e^{-ikx}) \sin mx \]
\[ + rB(\chi)e^{-2kx} \sin 2\pi y, \quad (A23) \]
where the asterisk denotes complex conjugation. \( k' = \beta/[2(\beta - \pi^2)] \) [see comment following \( (A15) \)] and \( k^2 = \beta - \pi^2 \). For \( \beta = 2\pi^2 \), \( k' = 1 \) and \( k = \pi \) and we shall restrict ourselves to this particular case. Substitution of \( (A23) \) into \( (A22) \) yields
\[ \nabla^2 \psi_x^2 + \beta \psi_y^2 + r \nabla^2 \psi^2 \]
\[ = \frac{1}{2} \sin(\pi y) e^{i \pi x - r x} \left[ \frac{dA}{dx} \right] (2 \pi^2 + 4ir \pi) \]
\[ + 4i \pi^4 r AB e^{-2 \pi x} \cos 2 \pi y \]
\[ + * + \pi^2 r \sin(2 \pi y) e^{-2 \pi x} |A|^2 \]
\[ + 2 \pi^2 r \frac{dB}{d\chi} e^{-2 \pi x} \sin 2 \pi y, \quad (A24) \]

where we assume that \( r \approx O(\epsilon) \) and neglect terms which are \( O(r) \) and the asterisk represents the complex conjugate of the first term on the right-hand side.

We seek a solution for the \( O(\epsilon^2) \) problem which is uniformly valid for \( x = O(\epsilon^{-1}) \). Consequently, the secularity producing terms appearing on the right-hand side of (A24) must be removed. This leads to two equations governing the spatial dependence of \( A \) and \( B \)

\[ \frac{dA}{d\chi} = i \pi^2 r \exp(-2r \chi/\epsilon) B, \quad (A25) \]
\[ \frac{dB}{d\chi} = -\frac{1}{2} \pi |A|^2, \quad (A26) \]

from which it follows that

\[ |A|^2 = |A(\chi_0)|^2, \quad (A27) \]
\[ B = B(\chi_0) - \frac{1}{2} \pi |A(\chi_0)|^2 \chi, \quad (A28) \]
\[ \frac{d\theta}{d\chi} = \pi^2 r (B(\chi_0) - \frac{1}{2} \pi |A(\chi_0)|^2 \chi) \]
\[ \times \exp(-2r \chi/\epsilon), \quad (A29) \]

where \( \theta \) is the phase of \( A \). The solution for the perturbation streamfunction is

\[ \psi = \epsilon |A(x_0)| \cos(\pi x + \theta) e^{-r x} \sin \pi y \]
\[ + \epsilon r [B(x_0) - \frac{1}{2} \pi |A(x_0)|^2 \epsilon x] e^{-2 \pi x} \sin 2 \pi y, \quad (A30) \]

which in the limit \( r \to 0 \) yields

\[ \psi = \epsilon |A(x_0)| \cos(\pi x + \theta) \sin \pi y. \quad (A31) \]

Since the forced solution of the \( O(\epsilon^2) \) problem is \( O(r) \) it follows that in the limit of \( r \to 0 \) the solution of the linear problem is an exact solution of the nonlinear equation. Q.E.D.

So far we have refrained from discussing the case when \( \beta = \pi^2 \). The far-field viscous solutions (A6) and (A7) suggest \( x \)-independent response. The viscous solutions (A13) and (A15) are ambiguous since we have to consider the double limit \( (\beta - \pi^2) \to 0, \quad r \to 0 \). To overcome the difficulties we have to let \( \beta = \pi^2 \) in (A1) and to consider the far-field solution. It is not difficult to show that the inviscid solution is unbounded downstream and, consequently, is inconsistent with the boundary conditions (A2). The viscous solution decays downstream but the decay rate is \( O(r) \) and there is uncertainty in estimating the amplitude of the perturbation. The difficulties associated with this case are removed once the complete solution is obtained. For \( \beta = \pi^2 \) and in the limit of \( r \to 0 \) the wavenumber for the stationary wave is \( k = 0 \). The group velocity associated with this wavenumber is zero. This implies that the wave energy is not radiated to infinity with a consequence of an amplitude buildup which can only be checked by friction. We have here a case of resonance. This problem is investigated by Merkine (1980) and the analysis shows that the amplitude of the perturbation is \( \approx O(\epsilon^{1/3}) \).

Fig. 2.1 depicts qualitatively the full streamlines, i.e., the unperturbed current plus the perturbation field in accord with our earlier discussion of the far field solution. The amplitude of the perturbation field has been exaggerated in order to exhibit the possible flow-field behavior in the presence of strong nonlinearity. Note the tendency for omega-type blocking for \( \beta > 2 \pi^2 \). The near-field is inferred from the energy equation (Section 4, positive generation of wave kinetic energy) or from potential vorticity conservation.

**APPENDIX B**

**Numerical Scheme**

We designed a simple scheme which conserves potential enstrophy and allows a straightforward implementation of most boundary conditions. We used a staggered grid with the prognostic variable \( \eta' \) (potential vorticity) defined at the center of the grid, and the streamfunction \( \psi \) at the corners (Fig. B1).

Using standard finite-differences nomenclature, i.e.,

\[ \delta x f_j = (f_{j+1/2} - f_{j-1/2})/\Delta x, \]
\[ f_{j+1/2} = (f_{j+1/2} + f_{j-1/2})/2, \]

and similarly for \( y \), the potential vorticity equation is

\[ \partial \eta'/\partial t = -\delta x [(U + u')\eta'], \quad (B1) \]
where $\eta_0 = fh/D$ is the orographic component of the potential vorticity, and the variation of planetary vorticity is not included in the definition of $\eta'$. The relative vorticity $\zeta'$ is defined at the corners of the grid:

$$\zeta' = (\overline{\eta - \eta_0})^x^y$$

(B2)

which allows the determination of the streamfunction and velocity components from

$$(\delta_x^2 + \delta_y^2)\psi' = \zeta'; \ u' = -\delta_y \psi'; \ v' = \delta_x \psi'. \quad (B3)$$

The boundary conditions are $\psi' = 0$ at all boundaries, $\eta' = 0$ at the upstream boundary and $\delta_x \eta' = 0$ at the exit.

We did experiments varying the distance between the mountain and the downstream boundary and verified that the solution in the interior is not influenced by the boundary. We used an $N$-cycle time difference scheme (Lorenz, 1971) with $N = 4$.

REFERENCES


———, 1950b: Blocking action in the middle troposphere and its effect upon regional climate. II. The climatology of blocking action. Tellus, 2, 275–301.


