Potential Vorticity Sources and Sinks in a Quasi-geostrophic Ocean:
beyond Western Boundary Currents

M. SUSAN LOZIER AND STEPHEN C. RISER

School of Oceanography, University of Washington, Seattle, Washington

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ABSTRACT

Potential vorticity dynamics for a quasi-geostrophic eddy-resolving general circulation model (EGCM) are studied in order to determine the effects of mesoscale variability on the potential vorticity distribution of a wind-driven ocean. The study employs both Eulerian and Lagrangian analyses in the effort to describe the potential vorticity gain/loss cycle along the path of a particle. While the mean wind stress curl is the dominant potential vorticity source for the interior of the upper layer, a redistribution of eddy potential vorticity creates sources of potential vorticity for the multiple gyres in the lower layers. This redistribution is a result of the local generation of eddies via baroclinic instabilities. These eddies are advected by the western boundary current into the midlatitude jet where they are responsible for a cross-gyre potential vorticity exchange. This exchange is concentrated at the entrance to the eastward jets where northward and southward boundary currents converge. From a Lagrangian viewpoint the vorticity exchange is accomplished via dissipative meandering rather than particle exchange across gyre fronts.

1. Introduction

In a previous work (Lozier and Riser 1989, hereafter LR) the authors examined the role of boundary currents in the vorticity cycle of a three-layer, wind-driven, eddy-resolving general circulation model (EGCM). Dissipative and inertial boundary currents were found in each of the three layers, with the former providing a sink for the potential vorticity gained in the ocean interior. However, within a double gyre ocean the boundaries are not the sole sink of potential vorticity since eddy fluxes can transfer potential vorticity from one oppositely-driven gyre to the next (Harrison and Holland 1981; Holland and Rhines 1980). In a study of a barotropic ocean model, Marshall (1984) restricted all vorticity loss to this transfer between counterrotating gyres through the use of unconventional boundary conditions that excluded any vorticity dissipation at the western boundary. Marshall’s study and Stommel’s (1948) study of western intensification constitute the two limiting cases of partitioning between potential vorticity loss via dissipative boundary currents and potential vorticity loss via cross-gyre transfer. The ratio of these two sinks, dependent upon the relative strength of inertial to frictional forces, is model and parameter dependent. For the representative model ocean chosen for this study both sinks exist. While the need for these sinks in a barotropic ocean is determined by the potential vorticity source supplied by the wind curl in the Sverdrup interior, the question arises as to the need for these sinks in the lower layers of a baroclinic ocean where there is no direct interior potential vorticity source.

In this paper the interior source distribution of potential vorticity is examined for all three layers, and the cross-gyre exchange is examined beyond the Eulerian measurement of a time-averaged potential vorticity flux that was shown in LR. The mechanism of cross-gyre exchange pertains not only to the midlatitude jet that exists in all three layers but also to the shared gyre fronts of the multiple gyres in the lower two layers. The intent of this study, in conjunction with the previous study, is to examine the nature of the potential vorticity gain/loss cycle for a wind-driven ocean model with mesoscale variability and to explore the manifestations of these mechanisms in the Lagrangian frame.

An extensive analysis of Eulerian potential vorticity dynamics was performed for a similar, but two-layer, EGCM by Holland and Rhines (1980, hereafter HR). In an effort to understand the dominant potential vorticity mechanism at various spatial scales, pointwise, gyre, and basin balances were examined. Alternatively, a Lagrangian path represents an integration of the spatial scales of the vorticity balance, and thus affords a different perspective on the dynamics. This study is an attempt to complement HR’s earlier Eulerian analysis with such a Lagrangian analysis.

2. Model equations and Lagrangian paths

The model equations, a schematic of the model’s vertical structure, and the parameter choices for the
realization discussed in this paper are given in LR. 
Briefly, the model dimensions are 4000 km × 4000 
km (L × L) with mean layer thicknesses of $H_1 = 300$ 
m, $H_2 = 700$ m and $H_3 = 4000$ m. The surface layer 
is driven by a sinusoidal wind stress that creates an 
anticyclonic (cyclonic) gyre in the southern (northern) 
half of the basin. Lateral dissipation is parameterized 
as Laplacian friction and bottom friction is parameterized 
as Rayleigh friction. For each layer $i$ ($i = 1, 3$) 
the quasi-geostrophic streamfunction is represented by 
$\Psi$, with the horizontal velocity field $(u_i, v_i)$ given by 
$u_i = (-\psi_{yi}, \psi_{xi})$. Each layer is governed by 
\[
\frac{Dq_i}{Dt} = F_i + D_i
\]
(1)
where
\[
\frac{D}{Dt} = \frac{\partial}{\partial t} + J(\psi_i,)
\]
with $J$ representing the Jacobian operator, and $F_i$ and 
$D_i$ the appropriate forcing and dissipation for each 
layer. The potential vorticity is designated as $q_i$ and is 
comprised of three components: planetary vorticity, 
relative vorticity and stretching vorticity, designated as 
$f$, $\xi$ and $\eta$, respectively. The planetary vorticity, $f$, is 
approximated by $f_0 + \beta(y - L/2)$, where $L$ is the basin’s 
north–south and east–west extent, and the stretching 
vorticity terms are given by
\[
\eta_1 = \frac{(\psi_2 - \psi_1)f_0^2}{g'_{12}H_1},
\]
(2a)
\[
\eta_2 = -\frac{(\psi_2 - \psi_1)f_0^2}{g'_{12}H_2} + \frac{(\psi_3 - \psi_2)f_0^2}{g'_{23}H_2},
\]
(2b)
\[
\eta_3 = \frac{(\psi_3 - \psi_2)f_0^2}{g'_{23}H_3},
\]
(2c)
where $g'_{12}$ and $g'_{23}$ are the reduced gravity values at the 
first and second interfaces, respectively.

Using Laplacian friction, the imposed lateral 
boundary conditions are limited to
\[
\zeta_i = 0,
\]
\[
u_i \cdot n = 0,
\]
where $n$ is the unit vector normal to the boundary. In 
addition to the quasi-geostrophic potential vorticity 
equations, there are continuity equations that govern 
the deviation of the two interfaces from their equilibrium 
position. These are given by
\[
\frac{Dh_{12}}{Dt} = \frac{\partial h_{12}}{\partial t} + J(\psi_{12}, h_{12}) = w_{12},
\]
(3a)
\[
\frac{Dh_{23}}{Dt} = \frac{\partial h_{23}}{\partial t} + J(\psi_{23}, h_{23}) = w_{23},
\]
(3b)
where $\psi_{12}$ and $\psi_{23}$ are interfacial streamfunctions 
evaluated from $\psi_i$ ($i = 1, 3$) via linear interpolation, $h_{12}$
and $h_{23}$ are the positive deviations of the interface 
heights, and the interfacial vertical velocities are given 
by $w_{12}$ and $w_{23}$. As the bottom is flat and the surface 
is rigid, the vertical velocities vanish at $z = 0$ and $z = -H$. In order to satisfy continuity, such that \[
\int w_{12}dxdy = 0 \quad \text{and} \quad \int w_{23}dxdy = 0, \]
the conditions $\partial h_{12}/\partial t \int h_{12}dxdy = 0$ and $\partial h_{23}/\partial t \int h_{23}dxdy = 0$ are 
imposed.

From an initial state of rest the quasi-geostrophic 
layered equations have been integrated in time until a 
statistical steady state was reached. The time-averaged 
streamfunction and potential vorticity fields for the 
numerical ocean are shown in Fig. 1 for each of the 
three layers. The averaged fields represent means over 
3600 days of Eulerian data, subsampled at 2-day 
intervals. In the text to follow an overbar will designate 
such a time-averaged value and perturbations from the 
time average will be designated with a prime. While the $\psi$ 
pattern of the upper layer clearly reflects the imposed 
mean wind forcing, the $\bar{\psi}$ pattern of the lower 
two layers, characterized by multiple counterrotating 
gyres (refer to LR Fig. 4), reflects eddy-induced 
circulation. The nature of these multiple gyres and their 
effect on the interior potential vorticity distribution will 
be discussed in section 3. Of note in the $\bar{\psi}$ fields are the 
two regions of reversal with depth of the meridional $\bar{\psi}$ 
gradient. In the upper layer a southward gradient of $\bar{\psi}$ 
that appears directly below the midlatitude jet is a 
familiar feature of strong inertial recirculation that has 
been noted in previous EGCMs. Of particular interest in 
the upper layer of this model ocean is the additional 
southward gradient of $\bar{\psi}$ extending from the southern 
wall to midgyre, where the wind stress curl and the 
mean interface deformation are maximum. In both of 
these regions the mean meridional gradient of $q$ is 
dominated by a strong southward height gradient that 
reverses the influence of $\beta$. These upper layer equatorward $\bar{\psi}$ gradients contrast the poleward gradients found 
in the layers below. As will be discussed, it is believed that 
baroclinic instabilities in these regions create the 
pattern of multiple gyres in the lower layers.

Within a statistically steady ocean clusters of 16 particles 
(“floats”) were released at 26 locations throughout 
each layer of the model ocean with the majority of 
the locations in the southern basin. Depending upon 
the area of study, one of three different deployment 
schemes, each with 26 clusters, was used. Thus, there 
are a total of 78 initial cluster locations. For simplicity, 
Fig. 2 gives the initial location of only those clusters 
referred to within this text. Floats from each release 
were tracked for an average of 182 days and then 
reinitialized at their original launch position for the next 
release. Twelve serial releases of a single deployment 
scheme were used in this study with six releases of 
the remaining two deployment schemes. Each release rendered 
75 712 float-days per layer. Individual paths, 
representative of the group, have been chosen to illustrate 
Lagrangian behavior for this study. Additionally,
the change in $q$ and its components (mean, eddy and/or instantaneous) along these particle paths creates a Lagrangian signature that is used to illustrate the local dynamics. The reader is referred to LR for details of the Lagrangian path construction and the computation of changes in potential vorticity and its components along a particle path.

3. The ocean interior

In a steady, wind-driven ocean the interior flow, far removed from western boundary currents and strong inertial recirculations, is expected to behave according to Sverdrup dynamics. Although spatial integration over the heart of the model interior ocean ($500 \leq y$
forcing. The westward propagation of these waves is disrupted at approximately 3500 km, where a nonlinear eddy field with spatial scales on the order of 100 km appears. Possible sources for these eddies west of 3500 km are direct advection from the jet and/or local generation via instabilities. The latter will be shown to be

\[ \beta (v_1 H_1 + v_2 H_2 + v_3 H_3) = \text{curl} \tau, \]

a Sverdrup balance holds locally (to within 2%) only in the far eastern portion of the model ocean. As seen in Fig. 3, from approximately \( x = 3500 \) to 4000 km, the wind stress curl, which completely dominates the nonconservative forcing field, essentially balances the \( \beta v_1 \) field; there is no mean flow in the second and third layers near the eastern wall. To the west of this demarcation, which is the approximate location of the midlatitude jet exit, eddy stresses are important to the time mean flow of the lower layers. The flow fields west of 3500 km in the lower layers are characterized by "turbulent" Sverdrup balances (i.e., \( \bar{u} \cdot \nabla q = -\nabla \cdot \bar{u} q \)), as described by Rhines and Holland (1979). The regions to the east and west of 3500 km also differ in their temporal variability. Barotropic waves fill the basin to the east of 3500 km and their spatial and temporal scales of 500 km and 45 days are in agreement with the Rossby wave dispersion relation.

The time variability in this part of the model ocean, where there is only mean forcing and an absence of instabilities, is attributed to the radiation of barotropic energy from the unstable jet exit, such as suggested in a study by Harrison and Robinson (1979), where an oscillating northern boundary was used to simulate strong current variability. Additionally, in a study by Haidvogel and Rhines (1983), localized time-dependent forcing (a "plunger") was shown to produce Rossby waves in regions far removed from the local

**Fig. 2.** Initial float locations within all layers. Each letter represents 16 floats either spread evenly across the bracketed line or arranged in a \( 4 \times 4 \) cluster. The floats are spaced 20 km apart except for cluster E where the spacing is 10 km.

**Fig. 3.** Potential vorticity balance for the far interior of layer 1. Contour interval scaled by \( 10^{-12} \text{ s}^{-2} \). (a) curl \( \tau / H_1 + \nu \nabla^2 \zeta \); (b) \( \beta v_1 \).
of dominant importance to the gyre dynamics of the lower layers.

A distinction between the two interior regimes also appears in the Lagrangian frame. The vorticity balances for floats seeded in the far interior (east of 3500 km) yield a picture of mean surface-limited Sverdrup dynamics in conjunction with strong variability. The trajectory and accompanying Lagrangian signature of an upper layer float launched in this region from Cluster K (Figs. 4a–b) show a wavelike motion superimposed on a slow southward drift. Since dissipation is negligible in this signature, as are the changes in both $\vec{\eta}$ and $\vec{\zeta}$, only the changes in the eddy vorticity components and the integrated wind forcing are shown. For this particular signature the latter is equivalent to $q - q_0$. As seen from the Lagrangian signature, the wind forcing is responsible for the particle's net southward drift, which results in an uncompensated loss of planetary vorticity. The shorter variations in $f$ are compensated by the eddy relative vorticity, consistent with a large-scale barotropic planetary wave having a period of 45 days. The same wave signature is evident from the potential vorticity balance for a Cluster K float in layer 3 that was launched at the same time and locale as the layer 1 float described above (Figs. 4c–d). However, as expected, there is no net displacement for this float. The periodic change in $q$ is due to the dominance of bottom friction (which is essentially $-c\vec{\xi}$ in this region) over lateral friction. Floats in this layer and the middle layer oscillate on the beta plane and experience no net change in their $q$ value.

While all Lagrangian signatures in the far interior are virtually the same, the middle interior (west of 3500 km) Lagrangian signatures are characterized by temporal and spatial variability. They do, however, share some general characteristics. These are shown by a float that drifts southwestward into the middle interior from its launch location, Cluster J, within the far interior.

Fig. 4. Lagrangian potential vorticity signature for floats released within the far interior from Cluster K. (a) Layer 1 float trajectory superimposed on $\vec{q}$ field with contour interval scaled by $100/f_0$ (origin designated by "O" and "+" marks every 10 days); (b) Changes from initial vorticity values; (c–d) As for (a–b) but for a layer 3 float.
(Fig. 5). This float signature also serves to illustrate that the transition between a classical Sverdrup interior and an eddy-dominated interior is abrupt in the Lagrangian frame as well. A barotropic wave signature is observed while the float is in the far interior; however, at day 120 the float crosses into the middle interior (at approximately 3500 km) and there is an immediate change in its Lagrangian potential vorticity balance. The most notable change is the introduction of higher frequency motions that are characterized by a negative correlation between the changes in eddy stretching vorticity and eddy relative vorticity. Such baroclinic features, with varying temporal scales, are present in the signatures in all three layers. Additionally, after the crossover, dissipation becomes important to the $q$ balance. While the wind forcing causes a persistent loss (in the subtropical gyre) of $q$, dissipation, which can be as large locally as the wind stress curl, can cause the loss to be intermittent. (In this particular signature dissipation is of the same sign as the wind stress curl.) Although interior $q$ sources are anticipated in the lower layers (since western boundary layers provide a dissipative sink) unambiguous changes in $q$ along particle paths in the middle interior were not found. This difficulty is attributed to the small scale temporal and spatial variability for $q$ in this region. As will be shown, a $q$ interior source may be inferred from an Eulerian statistical mean of eddy $q$ fluxes. However, it is helpful to first examine the nature of the multiple gyre pattern in the lower layers. The appearance of such gyres in the lower layers of EGCs was first explained by Rhines and Holland (1979) in terms of a vorticity-transport model that depends upon the parameterization of the explicit eddy field in terms of a Lagrangian diffusivity. Basically, the mean abyssal flow field was explained in terms of the observed large scale $q$ field and the measured distribution of the calculated Lagrangian diffusivity. In this work an alternative, but compatible, viewpoint is presented that makes use of the eddy field explicitly.

As seen in Fig. 1 the mean $q$ fields exhibit two regions where there is a reversal of the meridional $\bar{q}$ gradient, thus creating a favorable environment for baroclinic instability. Both of these regions are characterized by strong downgradient eddy height fluxes (I.R.), which, in a quasi-geostrophic model, indicate energy conversion via baroclinic instabilities. The meridional distribution of the zonally integrated eddy height fluxes, $\int_0^L \bar{v}'_1 \bar{h}_{12} dx$ and $\int_0^L \bar{v}'_2 \bar{h}_{23} dx$ (Fig. 6a), shows that the flux maxima for each interface coincide with the regions of instability. The broader span of large flux values in the southern portion of the basin reflects the broader $\bar{q}$, reversal region noted in Fig. 1. Convergences and divergences of the eddy height flux straddle the maxima, thus contributing to mean vertical velocities at the interfaces as shown by an expansion of Eqs. (3a-b) into mean and eddy components:

$$\bar{\omega}_{12} = \mathbf{\nabla} \cdot \bar{\mathbf{u}}''_{12} \bar{h}_{12} + \bar{\mathbf{u}}'_{12} \cdot \mathbf{\nabla} \bar{h}_{12} \quad \text{(4a)}$$

$$\bar{\omega}_{23} = \mathbf{\nabla} \cdot \bar{\mathbf{u}}''_{23} \bar{h}_{23} + \bar{\mathbf{u}}'_{23} \cdot \mathbf{\nabla} \bar{h}_{23}. \quad \text{(4b)}$$

Although the interfacial velocity is inconsequential for the surface layer’s mean circulation pattern since it is dominated by the surface Ekman velocity, it is instrumental to the lower two layers as seen by succinctly writing the spatially integrated interior (exclusive of boundary currents) vorticity balances for the lower two layers as
Fig. 6. (a) Zonally integrated meridional eddy height fluxes. (b) Meridional distribution of zonally integrated potential vorticity fluxes.
Layer 2: $\beta \tilde{v}_2 = f_0 (\tilde{w}_{12} - \tilde{w}_{23})/H_2$  \hspace{1cm} (5a)

Layer 3: $\beta \tilde{v}_3 = f_0 \tilde{w}_{23}/H_3$  \hspace{1cm} (5b)

The gyre structure of the lower layers is established by the mean interfacial vertical velocities which, according to Eqs. (4a–b), are comprised of both mean and eddy components. When the eddy component dominates, the gyre structure is independent of the overlying interface or gyre and the flow is considered to be driven by form drag.

Given the two eddy height flux maxima (Fig. 6a) one would expect four eddy-driven gyres to result in the southern basin based on opposing vertical velocities on either side of the maxima. However, for the middle layer, only the divergent eddy flux near the southern boundary is sufficient to dominate $\tilde{w}_{12}$ such that a strong local upwelling velocity results. Similarly, a divergent eddy flux in the bottom layer produces a positive, but weaker, vertical velocity, $\tilde{w}_{23}$. The resultant stretching ($\partial \tilde{w}/\partial z > 0$) leads to a convergent horizontal flow field and hence a cyclonic gyre borders the southern wall. Elsewhere the middle layer flow is dominated by the mean contribution to $\tilde{w}_{12}$. Only in the bottom layer are the eddy-induced vertical velocities completely dominant such that the $\tilde{v}$ pattern reflects the pattern of the instabilities via the demarcation of the cyclonic and anticyclonic gyres. The pattern of the two flux maxima produces alternating divergent and convergent fluxes that drive the two cyclonic and two anticyclonic gyres, respectively, in the southern basin. The generation of these deep multiple gyres is, of course, model and parameter dependent. Although there is some observational evidence of a deep countercurrent below the westward return flow of the North Atlantic subtropical gyre (Riser and Rossby 1983; Hogg 1983; Wunsch and Grant 1982; Reid 1978), there is little reason to believe that the gyre pattern in the deepest layer of this ocean model is representative. These gyres could be suppressed with topographic effects, changes in wind strength and structure, and also changes in the instability mechanisms. The deep gyres in this model reflect the contribution from eddies alone, as pointed out by Harrison (1982) in a study on deep mean flow within quasi-geostrophic EGCMs and other, more complex models.

These gyres are also separated by maximum eddy $q$ fluxes since downgradient eddy height fluxes are accompanied by downgradient eddy $q$ fluxes in the model interior (where $\eta$ dominates $\zeta$). An inspection of the meridional distributions of $H_1 \int_{L}^{L'} \tilde{v}/\tilde{q}_1 dh = -\sigma_1 \tilde{q}_1 dh$ (Fig. 6b) reveals the presence of strong downgradient $q$ fluxes in the regions of $\sigma_1$ reversals with depth. In the upper layer strong positive $\tilde{v}/\tilde{q}_1$ fluxes (at $y = 300$ km and $y = 1800$ km) are associated with the negative meridional $\tilde{q}$ gradient in these areas, while the positive meridional gradient of $\tilde{q}$ in the lower layers are associated with negative flux values. As expected from the presence of the homogenized pool of $q$ in the middle layer, a vorticity source for the middle layer appears at the southern instability region only. However, for the bottom and upper layer, a strong vorticity input is present at both instability locales. The presence of local maxima indicate that within a given latitudinal band, potential vorticity is being gained (lost) due to a net eddy adveactive flux into (out of) the area. Therefore, the eddy advective fluxes are sources of potential vorticity (positive or negative) for those regions which straddle the flux maxima. In the upper layer these internal sources are overshadowed by the strength of the vorticity source provided by the wind, but for the lower two layers the eddy advective fluxes are vital to the dynamics of the multiple gyres. For the lower layers the flux maxima separate counterrotating gyres, as seen from an inspection of the $\tilde{v}$ fields. In agreement with this division there is a corresponding minimum of $\int_{L}^{L'} \tilde{v} q_1 dh$ at these latitudes. It is the zonally integrated eddy $q$ flux that closes the potential vorticity balance for the individual gyres. The match between the flux maxima and the gyre boundaries simply results from the aforementioned correspondence between eddy height fluxes and eddy $q$ fluxes. These internal vorticity sources are simply the transfer of vorticity from one gyre to another via eddies. Therefore, adjacent gyres in the lower layers have opposing signs for their potential vorticity source. This is consistent with alternating northward- and southward-flowing western boundary currents at the western wall that have opposing signs for their potential vorticity sinks. In summary, the westward return flows, marked by instabilities, gain potential vorticity (negative or positive) as they head toward the western wall. Within the western boundary currents this vorticity gain is only partially dissipated (LR). The remainder is lost along the eastward jets of these gyres, as will be discussed in section 4.

4. Cross-gyre vorticity exchange

Within this EGCM a single, well-connected and well-defined jet is located at midlatitude, due to the anti-symmetry of the wind pattern. The jet, a seaward extension of the western boundary currents, extends several thousand kilometers into the interior and is characterized by strong lateral and vertical shear. The jet exhibits strong spatial and temporal variability; it is marked by large-scale meanders that propagate downstream and grow in amplitude until the jet loses its coherence. The downstream break-up of this jet has been attributed to both barotropic and baroclinic instabilities of instantaneous jet fields in a study by Haidvogel and Holland (1984) of a similar, but two-layer EGCM. In this section the consequences of the jet's variability for its potential vorticity dynamics are examined. A time-mean potential vorticity flux across midlatitude may result from a diffusive flux, and/or
an eddy advective flux. For this model realization the latter contribution to the total flux is 2 to 3 orders of magnitude larger than the dissipative flux, as was shown in LR. Although the eddy advective flux has been measured in the Eulerian frame the question remains concerning the mechanism of this midlatitude potential vorticity exchange: How is this transfer of $q$ across oppositely driven gyres achieved and how is it manifested in the Lagrangian frame?

a. Overshooting: The influence of the western boundary

Despite the antisymmetrical wind forcing, strong temporal variability distorts the instantaneous gyre

![Diagram](image-url)

**Fig. 7.** (left) Instantaneous streamfunction fields for the midlatitude jet. Contour intervals are scaled by $10^{-3} \text{m}^3\text{s}^{-1}$. (right) Instantaneous potential vorticity fields for the midlatitude jet. Contour intervals are scaled by $100/f_0$. 
front from the line of zero wind stress curl, which approximately divides the time mean gyres. Such variability creates an instantaneous midlatitude jet that is characterized by intermittent overshootings from the northern and southern western boundary currents. Shown in Fig. 7 are the instantaneous $\psi'$ fields during an overshooting from the south, which has traveled approximately 200 km beyond midlatitude. (The meandering is exaggerated in these fields due to the expansion of the meridional scale.) Of note in these fields is the spatial periodicity of the jet, the vertical coherence of the meandering pattern, and a sharp vertical shear. A time evolution of such overshootings is shown by a phase diagram of $\psi'(x, t)$ located at mid-latitude (Fig. 8). Attention is first focused at the jet entrance (near $x = 0$) where the northern and southern western boundary currents converge. At the beginning of this 360 day time series a weak overshooting from the north ($\psi' < 0$) is followed by two stronger southern overshootings ($\psi' > 0$) with an intermittent relaxation of the jet; a pattern that is repeated in the lower two layers for the same time series. Such change in conditions at the jet entrance causes an intermittency in the jet pattern's eastward propagation. As is evident in Fig. 8, the propagation (at approximately 10 km day$^{-1}$) occurs only in response to varying initial conditions at the jet entrance. When the initial condition is relatively constant the meander of the inertial jet resembles a stationary quasi-geostrophic wave. There is no temporal periodicity associated with this variability in the jet entrance, as measured by an autocorrelation of $\psi'$ with lag periods extending to five years. Instead, the changing initial conditions are a result of the aperiodic temporal variability in the western boundary currents. The asymmetry of the southern and northern basins' flow fields, introduced by mesoscale variability, causes the northern and southern western boundary currents to deposit eddies of varying strength into the jet entrance at varying times; thus causing aperiodic overshootings. These eddies originate from local generation in the unstable westward flows of the interior, and are advected into the western boundary current and, subsequently, into the jet entrance. A time evolution of an eddy being advected into the jet entrance is illustrated with contoured fields of the eddy potential vorticity, $q'$, in layer 2 of the northern basin (Fig. 9). Strong eddy activity is clustered at the northern edge of the homogenized area, where baroclinic instabilities generate midocean eddies. At $t = t_0$, a large pool of $q' > 0$ approaches the western boundary. This anomaly is subsequently advected down the western wall (with some attendant dissipation) and then injected into the midlatitude jet, where it causes a sharp increase in the eddy variability at the jet entrance from $t = t_1$ to $t = t_2$. Similarly, eddies enter the jet from the southern western boundary current. The appearance of these eddies in the boundary currents is shown by a phase diagram detailing $\zeta'(y, t)$ within the western boundary current. Figure 10 shows the path of variability along the length of the western boundary in the southern basin over a 360 day period. The appearance of strong eddies is aperiodic and their travel up the western boundary, at approximately 25–30 km day$^{-1}$, is due mainly to advection by the mean current. (Southward advection at the southern portion of the western boundary results from the presence of the cyclonic gyre at the southern wall in this middle layer.) Of particular interest in this phase diagram is the change in value of $\zeta'$ at the jet entrance. Negative perturbations within the dissipative layer reach the entrance from the south while positive perturbations reach the entrance from the north. This switch in $\zeta'$ values at the entrance simply reflects the alternating overshooting of the two western boundary currents as previously shown for the $\psi'$ fields at the jet entrance. An isolated eddy with $\psi' > 0$ will have $\zeta' < 0$ due to its curvature, and will also be characterized by $\nabla^2 \zeta' > 0$ next to the western wall where $\zeta' = 0$ at $x = 0$. Stronger dissipative values within the boundary layer will accelerate the change in potential vorticity according to 

$$\frac{Dq}{Dt} = \nu \nabla^2 \zeta'. $$

As discussed in LR the change in $q$ within the western boundary is dominated by planetary vorticity changes both in the mean and instantaneously. Therefore, the eddy potential vorticity balance is predominantly

$$\beta \theta' = \nu \nabla^2 \zeta';$$

thus negative relative vorticity anomalies (which lead

![Fig. 8: Phase diagram of $\psi'(x, t)$ at $y = 2000$ km in layer 1. Contour interval is $30 \times 10^3$ m$^2$ s$^{-1}$.](image-url)
to positive dissipative anomalies) are associated with positive meridional velocity anomalies. A phase diagram of $\beta \psi'$ (Fig. 10, which corresponds to the same time series as shown for $\xi'$) illustrates this association between $\psi'$ and $\xi'$. Where $\psi'$ is positive (negative) the southern (northern) western boundary current overshoots and carries it negative (positive) $\xi'$. Under the constraints of quasi-geostrophy there is a strong positive correlation between $\xi'$ and $\eta'$ in the upper layer. Therefore, despite the change in $\psi'$ (hence $\psi'$ and $\xi'$) at the jet entrance the value of not only the eddy relative vorticity flux, but also the eddy potential vorticity flux, $\eta'\psi'$, remains negative. However, in the lower layers, where despite the strong vertical shear there is no flow reversal, $\xi'$ and $\eta'$ are negatively correlated (e.g., when the upper layer is stretched, the lower layers are squashed, yet the induced relative vorticity is of the same sign in all three layers). Because relative vorticity perturbations in the lower two layers generally outweigh those of stretching vorticity in this boundary regime, $v_2^2q_2^2 < 0$ and $v_3^2q_3^2 < 0$, still result although much weaker relative to the upper layer eddy flux. Therefore, in all three layers there is a transfer of $q$ from the northern basin into the southern basin at the jet entrance. The dynamics of the jet entrance may also be understood by considering the Lagrangian $q$ gain/loss cycle. For a particle which has gained excess negative (positive) potential vorticity from the interior in the southern (northern) basin, the western boundary alone does not provide sufficient dissipation. The particle must then overshoot its equilibrium position at midlatitude for further dissipation, as discussed in section 4b.

b. Midlatitude vorticity exchange

The mechanism of the midlatitude vorticity flux may be explained from a Lagrangian viewpoint in two ways: 1) particles meander from one gyre to the next return-
Fig. 10. (upper) Phase diagram of $\beta_2(y, t)$ within the western boundary layer. Contour interval is $2 \times 10^4 \beta_2/\alpha_2$ (lower) Phase diagram of $\beta_0(y, t)$ within the western boundary layer. Contour interval is $10^{-12} \text{s}^{-2}$. Values in excess of $|10^{-11} \text{s}^{-2}|$ are not contoured.

ing to their gyre of origin at the jet exit, and/or 2) particles permanently cross from one gyre to the next, carrying anomalous potential vorticities. In the first case dissipation is required in order that the vorticity carried back into the original gyre is changed in value, otherwise no $q$ transfer would occur. [As mentioned by Marshall (1984) a particle's meandering path allows time for more dissipation to act.] Ideally, these two choices could be discriminated by tagging particles in the western boundary currents, measuring the amount of vorticity carried into the adjacent gyre by those particles that escape permanently from their gyre of origin, and comparing that vorticity to the vorticity change of particles that return to their gyre of origin. Unfortunately, the unreliability of the $q$ estimates along float paths in the jet region, as discussed in LR, precludes such a determination. Nevertheless, the tracks alone are an invaluable aid in assessing the extent of cross-gyre communication and will be used along with Eulerian characteristics to study these two flux mechanisms.

1) DISSIPATIVE MEANDERING AS A MEANS OF VORTICITY EXCHANGE

This method for vorticity exchange must rely on the decay of $q$ along a particle’s path. Because the measured vorticity exchange is accomplished via an eddy advective flux it is the decay of $q'$ that is of interest. Such decay is explored with the use of a Lagrangian enstrophy, $q'^2$, equation that may be written as

$$\frac{1}{2} \frac{Dq'^2}{Dt} = -q' \frac{Dq}{Dt} + \overline{Dq'}$$

(6)

where for the upper two layers $D' = \nabla \cdot \nabla' \zeta'$, for the bottom layer $D' = \nabla \cdot \nabla' \zeta' - \zeta \zeta'$ and for all three layers $D' \gg D$. Additionally, $\overline{Dq'}$ is considered negligible in this region (HR). The first term in Eq. (6) represents the change in enstrophy along a particle’s path, the second term represents the change in enstrophy due to a particle’s meandering across the mean $q$ field, and the third term represents enstrophy dissipation. In their study, HR discussed the role of the enstrophy cascade in allowing particles, on average, to cross mean $q$ contours. Here, the specific interest is how the loss of enstrophy by individual particles relates to the structure of the cross-frontal exchange of $q$. For the case of a meandering particle returning to its initial $q$ value, it is the balance between the first and third terms that are of interest. As seen in Fig. 11, the mean enstrophy dissipation field, $\overline{Dq'}$, in the upper layer’s jet is negative everywhere with a threefold decrease in the absolute values from the jet entrance to the jet exit. The negative enstrophy dissipation values ensure the decay of enstrophy as a particle travels within the jet, so that even if a particle returns to its gyre of origin it has effectively transferred potential vorticity into the neighboring gyre. The mean enstrophy dissipation values are negative also in the lower layers but the $\overline{Dq'}$ field is essentially restricted to the jet entrance since the values outside of the jet entrance decrease by approximately two orders of magnitude.

This difference in the zonal extent of the $\overline{Dq'}$ field between the layers may be explained by examining the enstrophy decay time scale for the layers of the model. A simple scaling of the two pertinent terms in Eq. (6) leads to an estimate of the enstrophy decay time scale, $t_d$, of,
\[ L_d = \frac{q'\lambda^2}{\nu \bar{S}'^2} = \frac{(\eta' + \xi')\lambda^2}{\nu \bar{S}'}, \]

where \( \lambda \) is the length scale of the relative vorticity field. Because dissipation is parameterized as \( \nu \bar{S}'^2 \), rather than \( \nu \bar{S}'^2 q' \), the magnitude of \( q' \), relative to its components, plays a role in the determination of the decay time scale for enstrophy. Due to the positively and negatively correlated eddy \( q \) components in the upper and lower layers, respectively, there are large differences in the relative size of \( q' \) among the layers, and hence there are large differences in the layers' decay time scales. To illustrate these differences, attention is first restricted to the decay within the jet entrance where the correlations between \( \xi' \) and \( \eta' \) are strong and uniform. For the upper layer, where \( \xi' \approx \eta' \) and where the western boundary dissipative length scale, \( \lambda = 17 \) km, is appropriate (L.R.), the time scale, \( t_d \), is estimated to be 33 days. This estimate far exceeds the time a particle spends in the jet entrance (approximately 5 days). Using a characteristic velocity of 100 km day\(^{-1}\) the length scale for decay is estimated to be 3300 km. In the middle layer, where \( \xi' \approx -0.9\eta' \), the time scale is approximately 1.7 days. For a characteristic velocity
in the middle layer of 40 km day\(^{-1}\), a dissipative length scale of 68 km results. The values are approximate, but the zonal extent of the enstrophy field in layer 1 (Fig. 11c) versus layer 2 (Fig. 11d) is in agreement with the estimated decay length scales. According to these decay scales, the meandering in the lower two layers is viewed as a deep expression of the upper layer jet, and beyond the jet entrance it is not necessary for enstrophy dissipation. As evidence of these enstrophy decay differences, potential vorticity anomalies from the western boundary currents may be tracked downstream in the upper layer jet; but in the second layer such anomalies rapidly disappear outside of the jet entrance.

In the Eulerian frame the dissipation of enstrophy allows for a downgradient eddy \(q\) flux as seen from the Eulerian conservation statement for time-averaged enstrophy, \(\overline{u \cdot \nabla q^2 + u'q' \cdot \nabla q} = \overline{D'q}\).

The first term on the left side represents the advection of enstrophy and is the sum of \(\overline{u \cdot \nabla q^2}\) and \(\overline{u'q' \cdot \nabla q}\); the enstrophy advection by the mean flow and eddies, respectively. The second term on the left side represents the production (destruction) of enstrophy by the downgradient (upgradient) eddy potential vorticity flux across mean \(q\) gradients. Finally, the third term represents the enstrophy dissipation that allows for the enstrophy cascade to small scales. In the absence of strong advection the persistently negative enstrophy dissipative values would create a downgradient eddy \(q\) flux everywhere along the length of the jet. However, as seen from the \(u'q'\) vector plots in Fig. 12, the eddy \(q\) flux has alternating up- and downgradient values along the path of the jet, such that a bull’s eye pattern results. Such a bull’s eye pattern (in heat flux as well as \(q\) flux) was attributed by HR to the “adiabatic” wandering of the gyre boundary, whereby warm (or cold) water is alternately carried north and south. Indeed, the instantaneous jet’s meandering into regions of high and low enstrophy creates the alternating regions of enstrophy production and destruction; thus the dominant balance in Eq. (7) is between the two terms on the left side. However, enstrophy dissipation is needed to close the balance, especially near the western wall where its contribution is of the same order of magnitude as the contribution from advection and production. As shown by Marshall and Schutts (1981) in a study of oceanic and atmospheric geostrophic eddy fluxes, the nondiagonal part of the eddy \(q\) flux balances the advection of enstrophy while the divergent flux balances the enstrophy dissipation. The dominance of the rotational pattern for the eddy fluxes, shown in Fig. 12, reflects the strength of inertial to frictional forces within the jet. However, a zonal integration from the western to the eastern wall produces a dominant contribution from the divergent flux. The advective term, with its alternating transports, is essentially self-canceling, as noted previously by HR. Thus, it is the persistently negative enstrophy dissipative values that allow for the transfer of \(q\) across the gyre front via a downgradient eddy \(q\) flux. Although this transfer is strongest at the western wall (where dissipative values are strong due to the imposed boundary condition and where the \(q\) anomaly is also largest) it extends the length of the jet in the upper layer.

The eddy \(q\) flux pattern for the middle layer (Fig. 12b) is in sharp contrast to the flux pattern for the upper layer. Although there is a strong downgradient flux within the jet entrance, there is a two-order of magnitude decrease in values downstream. This is in agreement with the enstrophy decay length scales for the middle layer estimated earlier. Potential vorticity anomalies are entirely dissipated within the jet entrance of the lower layers and do not appear downstream. As mentioned earlier, such rapid dissipation is due to the relative smallness of \(q'\) in the jet entrance. Likewise, the eddy \(q\) flux is small relative to the eddy \(\eta\) and eddy \(\xi\) fluxes, as shown in Fig. 12. The counterrotation of these fluxes, due to the negative correlation between \(\eta\) and \(\xi\), effectively cancels an eddy \(q\) flux. The corotation of the \(\eta\) and \(\xi\) fluxes (Fig. 12) in the upper layer produces a sizable eddy \(q\) flux that extends the length of an instantaneous jet.

2) PARTICLE EXCHANGE AS A MEANS OF VORTICITY EXCHANGE

Particle exchange across fronts, such as the Gulf Stream, has been widely documented in observational float programs (Owens 1984; Shaw and Rossby 1984; Rossby et al. 1985). Additionally, float exchange across the midlatitude jet is present within this model ocean, with the degree of exchange temporally and spatially variable. This section addresses the question as to what relationship there is between observed float exchange and vorticity transfer. Does particle exchange contribute to the eddy potential vorticity flux across midlatitude, or does dissipative meandering suffice?

As a first step the tracks of floats within the jet were analyzed in order to assess 1) which floats escape to a neighboring gyre and 2) the location of float escape. From model runs with floats seeded north and south of the jet entrance, a particle’s duration within the jet is found to be strongly dependent on its initial launch position. This dependency is exemplified by 45-day trajectories of Cluster E floats launched in the northern gyre of the upper layer (Fig. 13). The 16 floats in this cluster were seeded 10 km apart at 400 km north of midlatitude. The floats spanned the width from \(x = 5\) to 155 km, thereby including both the dissipative and the inertial boundary layers which extend 20 km and 70 km from the western wall, respectively (LR). Float 1, which was launched initially within the \(q\) boundary layer at the western wall, is the only float that travels the full extent of the jet. This float, at the jet exit, escapes
to the southern gyre. As initial distance from the wall increases, the floats are seen to leave the stream earlier. These floats are recirculated back into their gyre of origin: the northern gyre in this case. Other clusters and other releases of floats launched at the jet entrance show the same general characteristics, as well as floats from the lower two layers.

Floats launched directly within the jet at various locations further characterize particle behavior within the jet. Clusters A, B, C and D are located at $x = 5, 100, 1400$ and $2600$ km, respectively within each layer. The 16 floats within each cluster are initially aligned N–S from $y = 2000$ km to $2300$ km, so that the floats initially span a full half-width of the instantaneous as
Fig. 13. Trajectories from one release of Cluster E in layer 1 with path length of 45 days. Distance from the western wall at launch increases with ascending order of float number.

well as the time-mean jet. The 45, 90 and 180 day trajectories for the first, second and third layers, respectively, are shown in Fig. 14. The trajectories show varying degrees of float dispersion with depth and with initial horizontal displacement from the western wall. Cluster A is unique in that all floats were launched within the thin dissipative western boundary layer. All 16 floats in the upper layer remain in the stream until the stream’s break-up, at which point the floats enter the southern gyre. In the middle layer, floats are seen to escape to the north and south, but only at the jet exit, which is westward of the exit region in the upper layer. This apparent decrease in jet penetration with depth is in agreement with a study by Holland and Schmitz (1985), which showed the sensitivity of jet length to inertial strength. The relative weakness of the jet in the lowest layer is evident from the trajectories of the Cluster A floats in that layer. Also of note for Cluster A floats is that the eastward meander propagation is evident in the lower layers due to the smaller jet velocities there. In effect, the slower that particles travel the more their streaklines will diverge due to the meander propagation. Cluster B floats, which are launched just outside the jet entrance, are drawn away from the stream along the jet’s length and only those few imbedded in the jet’s center complete the journey. Floats launched within Cluster C show only weak remnants of the mean jet path in the upper two layers and strong eddy activity causes reentrainment of floats launched at the edge of the jet. Finally, for Cluster D, which has its initial location near the jet exit, there remains no trace of the jet and the float tracks are characterized by strong eddy activity only. From this analysis of float trajectories it is clear that 1) eddy variability within the jet increases rapidly downstream and with depth; 2) floats within the jet’s center, those drawn from the dissipative western boundary layer, have longer path lengths within the jet, and 3) floats leave the jet all along its length but actual escapes (to the southern gyre in this case) occur near the jet exit only.

Due to the time variability of the jet these float tracks deviate from Eulerian streamlines, significantly so in the lowest layer. Regier and Stommel (1979), in a study of float trajectories within kinematic flow fields, demonstrated the sensitivity of the difference between Lagrangian tracks and Eulerian streamlines to the differences in \( u \) and \( c \), where \( u \) is the mean flow and \( c \) is the eddy flow velocity. In an observational study of SOFAR floats within the Gulf Stream, Owens (1984) postulated that the reason deep floats leave the stream to a greater degree than upper-layer floats is the difference between \( u \) and the characteristic propagation speed of the meanders. (Owens defined the propagation speed as the speed in the change of meander size.) The trajectories in this model offer support for that suggestion. In the upper layer, especially at the jet’s center where \( u > c \), the particles essentially are at the jet exit before the meandering could significantly adjust their location within the stream. To these floats the jet is essentially a stationary wave. In the lowest layer, however, \( u \ll c \), and float dispersion is evident at the early stages of the jet development. While the kinematic argument regarding float escape in the stream explains why floats leave the stream, the dynamic constraint imposed by potential vorticity is next discussed in order to explain why floats leave the jet where they do.

To understand the effect of float escape on the potential vorticity transfer, the instantaneous \( q \) fields are analyzed (Fig. 7). Of note in layer 1 is the narrowness of the jet defined by the \( q \) field compared to the jet defined by the instantaneous \( \psi \) field, which is also shown in Fig. 7. (Both the \( q \) and the \( \psi \) fields are from day 22 after the launch of the floats described above.) This thin \( q \) jet may be viewed as an extension of the dissipative boundary current alone, while the broader \( \psi \) jet encompasses the inertial boundary current also. Therefore, floats traveling within the jet’s center are embedded within a strong \( q \) gradient, which is an impediment to leaving the stream. Indeed, to cross such a \( q \) gradient during the time a particle is in the jet (\( \Delta q/\Delta t \)) requires dissipation values on the order of \( 10^{-10} \) \( \text{s}^{-2} \), which is an order of magnitude larger than the measured instantaneous values. The necessary dissipation would require a change on the order of 1000 \( \text{cm s}^{-1} \) over the width of the \( q \) jet (40 km). However, at the jet exit, where instabilities have created small-scale eddies, dissipative values increase, thereby facilitating escape from one gyre to another across a diminished \( q \) gradient. However, at this point in a float’s travel along the jet there is little enstrophy left to ex-
change (Fig. 11) due to the persistent decay along the jet's length. In fact, if the jet exit were dominated by this mixing there would be a strong negative $v q'$ signal in this region. However, $v q' > 0$ at the jet exit region (LR), reflecting the weakening enstrophy field along the path of the jet ($u \cdot v q'^2 < 0$) that forces an upgradient eddy $q$ flux.

Outside of the $q$ jet, but within the $\psi$ jet, floats are not restricted by dissipative constraints and they are free to leave the stream. Their exit, however, is restricted to their gyre of origin for the upper layer since otherwise they would have to cross the prohibitive $q$ gradient at the jet's center. Regardless of position relative to the jet's center, floats in the lower layers are not restricted by dissipative constraints as revealed by the lack of $q$ gradients in the instantaneous fields (Fig. 7). Thus, although there is ample float exchange beyond the jet entrance in these layers there is no accompanying vorticity exchange; it is largely restricted to the jet entrance. Where the $q$ gradients do occur in the profiles for the lower two layers there is no flow. This is the region of the far interior discussed earlier. As evidence of increasing escape with depth, of 4160 floats released within the southern gyre upper layer only 6% escaped to the northern gyre. This compares with 30% and 50% for the second and third layers, respectively.

As mentioned earlier, ideally the contribution of these two mechanisms, float escape and dissipative meandering, to the total $q$ transfer between the gyres could be quantified. However, lacking the particles' $q$ values in the jet, the Eulerian characteristics and the Lagrangian tracks must be used to infer the dominant mechanism. Because 1) vorticity exchange is concentrated at the jet entrance where dissipation is strong, 2) particle overshooting is prevalent at the jet entrance, 3) less than 1% of particles which overshoot from the western boundary currents leave the jet at the entrance, and 4) particles cross into neighboring gyres only where enstrophy values are low, it is concluded that the meandering mechanism rather than float escape is the dominant mechanism of vorticity transfer. Therefore, taking float exchange alone as a measure of vorticity transfer would lead to a serious underestimate. Improved estimates of $q$ following a particle within highly variable flows are needed before a quantification can be made. It should be noted that near the western wall, where the stretched grid creates a higher resolution, the $q$ estimates are valid. An estimated Lagrangian flux from particles $v'$ and $q'$ values as they passed the midlatitude yields values that agree to within 5% of those measured in the Eulerian frame. Additionally, as pointed out by HR, the breaking of meanders or waves (rings) contribute to the midlatitude $q$ flux. The effect of these rings has not been explicitly accounted for in this study (they are included in the eddy terms) although their contribution to the flux of dissolved ox-
ygen across the Gulf Stream has been estimated to be insignificant relative to an instantaneous diffusive flux (Bower et al. 1985).

Finally, although this discussion concerning cross-gyre vorticity exchange has exclusively concerned the eastward midlatitude jet, there are eastward jets separating the multiple gyres of the lower layers. The jet entrances of these eastward currents are also marked by strong eddy potential vorticity fluxes that result from the alternate overshooting of convergent eddy-containing boundary currents. The eddy $q$ flux serves the same purpose in the jet entrances of these eastward jets as it does for the midlatitude jet, namely to dissipate the vorticity gained in the interior.

5. Summary

Although each of the potential vorticity contributions, i.e., the boundaries, the generation of eddies in the ocean interior, the wind forcing, and the meandering of the eastward jets, has been discussed separately (boundary dynamics discussed in LR), their linkage is elucidated by long-term particle trajectories that also summarize the Lagrangian characteristics of

Fig. 15. Four-year spaghetti diagrams for layers 1, 2 and 3 ($0 \leq x \leq 4000$ km, $0 \leq y \leq 4000$ km).
the model circulation. Four-year “spaghetti” diagrams of six floats released at two locales within each layer of the model basin are shown in Fig. 15. (Although 16 floats were used for this diagnosis, six representative floats are shown for the sake of clarity.) The floats launched within the far interior of the upper layer (Cluster I) trace an outer dissipative track. After launch the floats move southward under the influence of the wind forcing. They are advedected westward along the southern wall, where their paths are highly variable, and then the floats are thread through the thin western dissipative boundary current and its seaward extension. The vorticity gained in the interior is primarily lost within the highly dissipative western boundary and within the jet entrance. The floats do not leave the eastward jet until the exit region as they are embedded in the thin q jet that restrains their escape. At the jet exit they are deposited at their origin in the far interior. Upper layer floats launched within the middle interior (Cluster F) trace an inner inertial track. After launch the floats are also carried up the western boundary but they flow within the broader inertial boundary current. Their travel along the eastward jet is marked by strong variability as the floats frequently leave and reenter the jet. The tight westward recirculation is evident from these float tracks.

Within the middle layer Cluster H was positioned at the edge of the homogenized q region. Trajectories launched from this cluster show two different behaviors. Some floats drifted into the far interior region where wave activity dominates and they do not escape from this region during the four years. The other floats outline the q homogenized area. They flow westward along the gyre front separating the anticyclonic and cyclonic gyres in the southern basin and are then advected into the dissipative boundary current and the eastward jet. Layer 2 floats launched from Cluster F trace the middle layer’s inner inertial track. The broader inertial western boundary current is evident as is the increased amount of escape along the jet’s length. Although the floats were launched in the southern basin, both the northern and southern basin’s westward recirculations are outlined by the float trajectories.

Finally, floats launched from Cluster I in the bottom layer illustrate the strong constraint placed on the flow by the background planetary vorticity field in the far interior. Without any source (or significant sink) the floats, subject to barotropic wave activity, are left to oscillate on the beta plane. In contrast, float trajectories that originate from Cluster G in layer 3 reveal a striking amount of structure. After four years of particle travel the two gyres nearest the zero wind stress curl line in both the northern and southern basins are outlined. Accordingly, both northward- and southward-flowing boundary currents are evident in this spaghetti diagram. The floats here clearly show the greater extent of cross-gyre float exchange in this lowest layer.

In conclusion, this analysis of a quasi-geostrophic ocean model has elucidated how mesoscale variability affects the potential vorticity distribution of surface and deep layers. In particular, it was shown that while the potential vorticity source for the upper layer gyres is dominated by the wind forcing, the lower layer gyres receive their potential vorticity interior source from eddy advective fluxes that transfer q across gyre fronts separating counterrotating gyres. These eddy q fluxes, resulting from baroclinic instabilities in the model ocean interior, create a redistribution of potential vorticity that is vital to the dynamics of each gyre in the lower layers. Also, it was shown that these interior-generated eddies cause overshootings of the western boundary currents, which leads to a transfer of q across midlatitude. This potential vorticity transfer, which is accomplished via dissipative meandering rather than float escape, is concentrated at the eddy-dominated jet entrance where enstrophy and dissipative values are large.

Finally, it is noted that dissipation plays both an explicit and implicit role in the model dynamics. In the presence of a rigid boundary, dissipation is explicitly important since dissipative fluxes are a strong q sink. However, even where dissipative fluxes are orders of magnitude smaller than advective fluxes (such as in the interior and at midlatitude) dissipation is crucial to the ability of the eddy advective fluxes to transfer potential vorticity across mean q gradients. In this respect it plays an implicit role. Therefore, because the overall potential vorticity distribution is vitally dependent upon dissipation, a major constraint of this model, in addition to the quasi-geostrophic approximation, is regarded to be the subgrid scale parameterization. It should also be noted that the results presented here are subject to quantitative changes due to variations of the model parameters and forcing functions. However, it is believed that the qualitative nature of these results is robust for model runs that yield a representative ocean circulation.

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