On the Interaction of Surface Waves and Upper Ocean Turbulence

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ABSTRACT

The phase-averaged energy evolution for random surface waves interacting with oceanic turbulence is investigated. The change in wave energy balances the change in the production of turbulent kinetic energy (TKE). Outside the surface viscous layer and the bottom boundary layer the turbulent flux is not related to the wave-induced shear so that eddy viscosity parameterizations cannot be applied. Instead, it is assumed that the wave motion and the turbulent fluxes are not correlated on the scale of the wave period. Using a generalized Lagrangian average it is found that the mean wave-induced shears, despite zero vorticity, yield a production of TKE as if the Stokes drift shear were a mean flow shear. This result provides a new interpretation of a previous derivation from phase-averaged equations by McWilliams et al. It is found that the present source or sink of wave energy is smaller but is still on the order of the empirically adjusted functions used for the dissipation of swell energy in operational wave models, as well as observations of that phenomenon by Snodgrass et al.

1. Introduction

Swell is the least-well-predicted part of the wave spectrum (Rogers 2002), and often causes surprises for operations at sea and on the coast. Snodgrass et al. (1966) observed a surprising conservation of energy of surface waves with 15–20-s periods, generated by a storm in the Southern Ocean and propagating undisturbed from New Zealand to Alaska, across turbulent equatorial regions with easterly winds and strong currents. In addition to interaction with the airflow that tends to dissipate waves traveling faster than the wind or against the wind (e.g., Kudryavtsev and Makin 2004), it has been suspected for a long time that waves interact with oceanic turbulence. Phillips (1961) underlined the expected different types of interactions of waves and turbulence, from energy-conserving wave refraction and scattering over current patterns larger than or of the order of the wavelength (e.g., McKee 1996; Bal and Chou 2002), to a wave-induced straining of much smaller turbulent eddies possibly resulting in the weak dissipation of wave energy.

Today's wave-forecasting models use poorly constrained parameterizations for swell "dissipation" that are necessary to reproduce observations of wave heights (e.g., Tolman and Chalikov 1996; Tolman 2002). Such a low-frequency dissipation complements the loss of wave energy resulting from wave breaking that mainly affects waves with periods shorter than the peak period, and effects of slow or opposing winds. Based on recent observations the dominant waves are known to break infrequently (Banner et al. 2000), depending on their steepness, but it is unclear how large the direct loss of energy is resulting from wave breaking for waves with frequencies at and below the peak frequency.

Because most wave energy is not directly dissipated into heat (viscous dissipation is generally negligible),

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the loss of wave energy may be studied through the associated production of turbulent kinetic energy (TKE). There is no doubt that breaking waves should be the main source of TKE in the top few meters of the ocean (Agrawal et al. 1992; Craig and Banner 1994; Melville 1996). This production of TKE is believed to occur mostly through the large shear at the forward face of breaking waves (Longuet-Higgins and Turner 1974) and through the shear around rising air bubbles entrained at the time of breaking.

2. Average shear below nonbreaking waves

We shall focus on the effect of this preexisting turbulence on nonbreaking waves. In this context the shear-induced production of TKE per unit volume, resulting from the organized wave and mean current motions, is expressed by

$$P_s = \overline{\rho_w u_i' u_j' \frac{\partial u_i}{\partial x_j}}, \tag{1}$$

where ρ_w is the water density, and indices i and j refer to any of the three Cartesian coordinates x, y, or z, with implicit sums over repeated indices. Primes denote turbulent fluctuations with zero means, while the overbar is an average over the turbulent realizations. We write u_i for the ith component of the Reynolds-averaged velocity, typically the sum of mean currents and wave-induced velocities. Because of the importance of surface shears, surface-following coordinates will be used, and the generalized Lagrangian mean (GLM) is chosen for its generality (Andrews and McIntyre 1978), denoted by an overbar with an L superscript. The GLM average is essentially an average over the mean trajectory of water particles, including the wave-induced (Stokes) drift.

It must be noted that this separation of the velocity field into turbulence, wave, and mean flow motions gives the same role to wave and mean velocities, which makes it different from the separation of Kitaigorodskii and Lumley (1983) who, motivated by data analysis, grouped together wave and turbulent velocities because of their overlapping frequency range.

Within the very thin surface and bottom boundary layers (respectively a few millimeters and a few centimeters thick), the inflexion point in the wave-induced velocity profile allows instabilities. The vertical turbulent flux of horizontal momentum is then generally proportional to the velocity shear, including wave and current velocities. In the bottom boundary layer this proportionality is well parameterized with eddy viscosities (e.g., Grant and Madsen 1979; Marin 2004) and yields useful expressions for the transfers of wave energy and momentum to turbulence and the bottom (Longuet-Higgins 2005). However, away from the boundaries no such relationship exists and eddy viscosities are clearly inapplicable. Indeed, mixing parameterizations for upper ocean currents use values of the vertical eddy viscosity K_z on the order of 0.1 m^2 s^{-1}. Such values of K_z in this context are meant to represent large fluxes of momentum in regions of weak velocity gradients, but a K_z of this magnitude would damp waves in less than a few periods, if turbulence acted like viscous forces represented by the molecular kinematic viscosity (Jenkins 1989). This could be called the "sea of molasses paradox," illustrating the unfortunate consequences of applying eddy viscosity values that are valid for unstable shear flows in results only valid for molecular viscosity and stable wave motions.

To proceed further we shall assume that the triple velocity correlation in (1) can be approximated as

$$\overline{P_s}^L = \overline{\rho_w u_i' u_j'}^L \overline{\frac{\partial u_i}{\partial x_j}}^L. \tag{2}$$

This may follow from assuming that the turbulent properties are not correlated with the wave phase as discussed below. To estimate the mean shears, one can then use the general relationship between an Eulerian average $\overline{\phi}$ of any variable ϕ and its corresponding GLM value $\overline{\phi}^L$, valid to second order in the wave slope [Andrews and McIntyre 1978, their (2.27)],

$$\overline{\phi}^L = \overline{\phi} + \overline{\xi_j \frac{\partial \phi}{\partial x_j}} + \frac{1}{2}\overline{\xi_j \xi_k} \frac{\partial^2 \overline{\phi}}{\partial \xi_j \partial \xi_k}. \tag{3}$$

Using linear wave theory for a monochromatic wave of amplitude a and a period compatible with wind-generated waves, and taking axis 1 in the direction of wave propagation and axis 3 pointing vertically upward, one has the following wave-induced velocities to first order in the wave slope and for weak current shears (Phillips 1977; McWilliams et al. 2004),

$$\tilde{u}_1 = a\sigma F_{\mathrm{CS}} \cos(kx_1 - \omega t), \tag{4}$$

$$\tilde{u}_2 = 0, \quad \text{and} \tag{5}$$

$$\tilde{u}_3 = a\sigma F_{\mathrm{SS}} \sin(kx_1 - \omega t), \tag{6}$$

where k and ω are the wavenumber and angular frequency, and σ is determined by the linear dispersion relation $\sigma^2 = gk \tanh(kD)$, where g is the acceleration resulting from gravity and D is the mean water depth. We have used $F_{\mathrm{CS}} = \cosh(kz + kD)/\sinh(kD)$ and $F_{\mathrm{SS}} = \sinh(kz + kD)/\sinh(kD)$. The wave-induced particle displacements are, to first order,

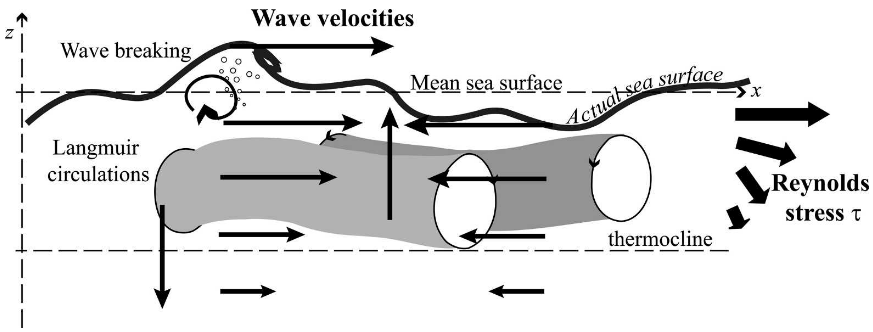

FIG. 1. Wave velocities (thin arrows) induced by long nonbreaking waves and the mixing processes in the upper ocean. Turbulent fluxes largely result from the presence of short breaking waves and may be carried in part by Langmuir circulations (gray "rolls"). In the limit (small slope for the long waves) in which these processes are not affected in the mean, and the mean shear induced by the long waves is not modified (small relative shear of the turbulent motions), the production of TKE resulting from the interaction of waves and turbulence is the turbulent momentum flux (thick arrows) times the volumetric mean of the wave-induced shears. That latter quantity is dominated by the shear under the wave crests.

$$\xi_1 = -aF_{\mathrm{CS}}\sin(kx_1 - \omega t), \tag{7}$$

$$\xi_2 = 0, \quad \text{and} \tag{8}$$

$$\xi_3 = aF_{\mathrm{SS}}\cos(kx_1 - \omega t). \tag{9}$$

Combining all of this into the generalized Lagrangian mean of (1), using (3), one sees that there are mean wave-induced shears that arise from correlations of shear and displacement,

$$\overline{\frac{\partial \tilde{u}}{\partial z}}^L = \overline{\xi_1 \frac{\partial^2 \tilde{u}}{\partial \xi_1 \partial \xi_3}} + \overline{\xi_3 \frac{\partial^2 \tilde{u}}{\partial \xi_3^2}} = \frac{a^2}{2} k^2 \sigma F_{\mathrm{CS}} F_{\mathrm{SS}} \tag{10}$$

and

$$\overline{\frac{\partial \tilde{w}}{\partial x}}^L = \overline{\xi_3 \frac{\partial^2 \tilde{w}}{\partial \xi_1 \partial \xi_3}} + \overline{\xi_1 \frac{\partial^2 \tilde{w}}{\partial \xi_1^2}} = \frac{a^2}{2} k^2 \sigma F_{\mathrm{CS}} F_{\mathrm{SS}}. \tag{11}$$

These expressions generalize to random waves because these second-order terms result from correlations of first-order quantities (see, e.g., Kenyon 1969 for a similar discussion). The two mean shears are each equal to half the vertical gradient of the Stokes drift U_s,

$$\overline{\frac{\partial \tilde{u}}{\partial z}}^L = \overline{\frac{\partial \tilde{w}}{\partial x}}^L = \frac{1}{2}\frac{\partial U_s}{\partial z}, \tag{12}$$

where $u = u_1$, $w = u_3$, $x = x_1$, and $z = x_3$. Therefore, the second-order wave-induced vorticity is obviously zero, because nonzero vorticity only occurs at a higher order, generated by boundary layers, the earth rotation, or current shears (see, e.g., Xu and Bowen 1994; White 1999; McWilliams et al. 2004). It may appear surprising that the average of $\partial w/\partial x$ is nonzero, which might be interpreted as leading to an infinite value of w as x goes to infinity. In fact, this nonzero mean corresponds to the fact that there is more water under the crests than under the troughs. Thus, the crests contribute more to the volume average (GLM preserves the volume at first order in the wave slope), and thus the first-order shear in the crests gives an indication of the second-order mean shears (Fig. 1).

In addition to the shear of the mean flow, these mean wave-induced shears produce TKE at the rate

$$P_{\mathrm{ws}} = \rho_w \overline{u_1' u_3'}^L \overline{\frac{\partial U_s}{\partial x_3}}, \tag{13}$$

as if the Stokes drift were a vertically sheared current.

It is interesting to discuss our assumptions, because this expression was previously derived by McWilliams et al. (1997) and Teixeira and Belcher (2002). The first authors arrived at this term by deriving the TKE equation from the Craik–Leibovich equations obtained from the Andrews and McIntyre (1978) GLM equation (Leibovich 1980; Holm 1996). This TKE production term appears as the result of the vertical gradient of a wave-induced radiation stress term (also called "Bernouilli head" in that context), and it also arises using the more general extension of the Craik–Leibovich equations given by McWilliams et al. (2004). Because these equations are phase averaged, it is natural that our assumption of no phase relationship between the turbulent fluxes and the wave motion yields the same expression.

Teixeira and Belcher (2002) instead considered the evolution of Reynolds stresses with the wave phase using rapid distortion theory, applicable to short gravity waves. With orthogonal curvilinear coordinates they found a modulation of $\overline{u_1' u_3'}$ at first order in the wave slope, and a second-order production of TKE equal to (13) arising from the correlation of this flux modulation with the oscillating wave shear. The equality of the two results suggests that their modulation of $\overline{u_1'u_3'}$ may be interpreted as a modulation of the Jacobian of their coordinate transformation, so that $\overline{u_1'u_3'}$ is constant in a volume-conserving coordinate transformation of the Cartesian coordinates. It should be noted that the Jacobian of the GLM transform does not fluctuate with the wave phase, and that the GLM coordinate transformation conserves the volume to first order in the wave slope (Jenkins 2004).

3. Application to swell dissipation

Now that we have seen that for short waves Teixeira and Belcher's (2002) results support the noncorrelation of $\overline{u_1'u_3'}$ with the wave phase that leads to (13), and that this expression may be obtained from the Eulerian mean equations of McWilliams et al. (2004), we may rely on (13) for application to all wave scales. As discussed, (13) is based on the hypothesis that one can neglect the mean effect of wave-induced modulation of the stress-carrying structures, typically vortices induced by wave breaking or Langmuir circulations. This hypothesis is expected to be well verified for swells of small amplitudes in the presence of a wind sea. In such a case, it is expected that the correlation of swell amplitude with the short wave breaking probability, and thus with the resulting horizontal distribution of $\boldsymbol{\tau} = (\overline{u'w'}, \overline{v'w'})$, only occurs at a higher order in the swell slope.

Based on the conservation of energy, a source of TKE is clearly a sink of wave energy and the rate of change of the wave spectrum $E(\mathbf{k})$ because only that effect may be expressed in the form of a source term,

$$\frac{dE(\mathbf{k})}{dt} = S_{\mathrm{turb}}(\mathbf{k}) = -\frac{2k\sigma E(\mathbf{k})}{g\sinh^2 kD}\int_{-D}^{0} \boldsymbol{\tau}\cdot\mathbf{k}\,\sinh(2kz + 2kD)\,dz, \tag{14}$$

where the wave vector $\mathbf{k}$ points in the direction of wave propagation. We shall now consider deep-water waves, for which

$$S_{\mathrm{turb}}(\mathbf{k}) = -\frac{4k\sigma E(\mathbf{k})}{g}\int_{-D}^{0} \boldsymbol{\tau}\cdot\mathbf{k}\,\exp(2kz)\,dz. \tag{15}$$

For relatively short waves $\exp(2kz)$ decreases away from the surface much faster than τ, so that the integral may be replaced by $\boldsymbol{\tau}(0)\cdot\mathbf{k}$. Equation (10) implies that turbulence damps waves propagating in the direction of the wind stress, while waves propagating against the wind would extract energy from turbulence. Yet, very close to the surface, wave breaking is a source of both TKE and momentum, concentrated over a distance that spans about 4% of the wavelength of breaking waves (Melville et al. 2002). Therefore, τ should increases from the surface on that scale. This should produce a "sheltering" of short waves from the Stokes shear dissipation, which will be neglected in the simple estimations shown here. This sheltering is expected to be significant for waves shorter than the dominantly breaking waves, namely, waves with phase speeds slower than about $5u_* \simeq U_{10}/6$, with u_* the air-side friction velocity and U_{10} the wind speed at 10-m height (Janssen et al. 2004). However, as suggested by anonymous reviewers and Kantha (2006), the orbital motion of long swells extends downward through a significant fraction of the mixed layer depth where τ rotates and decreases (Fig. 1). In general, the evaluation of S_{turb} requires the use of a mixed layer model, taking into account density stratification, in order to determine the vertical profile of τ. For deep mixed layers where the stratification may be neglected over the Stokes depth $1/2k$, τ varies on the vertical scale $\delta = \tau(0)^{1/2}/(4f)$, with f being the Coriolis parameter (Craig and Banner 1994).

To evaluate the magnitude of S_{turb} we shall use the approximation $\tau = \tau(0)\exp(z/\delta)$. This provides results within a factor 2 of those given by Kantha (2006). For a wind speed $U_{10} = 10\ \mathrm{m\ s^{-1}}$ the variation of τ reduces the depth-integrated dissipation by more than 50% for periods larger than 15 s. Calculations using the model of Craig and Banner (1994) suggest that the component of the stress perpendicular to the surface stress is always smaller than 25% of the surface stress but lies to the right of the surface stress in the Northern Hemisphere, possibly causing a slightly enhanced dissipation of waves propagating to the right of the wind.

Equation (15) implies that turbulence damps waves propagating in the direction of the wind stress, while waves propagating against the wind would extract energy from turbulence. That growth of opposing swells may be interpreted as a transfer of energy from the wind sea to the swell because the addition of an opposing swell reduces the mean shear because of the wind sea. In that case the total wave field loses energy as if it were a weaker wind sea with a smaller dissipation, the difference being pumped into the swell. One may further interpret the present result as the stretching or compression of Langmuir circulations (LCs) by the Stokes drift shear (McWilliams et al. 1997). On the oceanic scale, (15) yields results that suggest very dif-

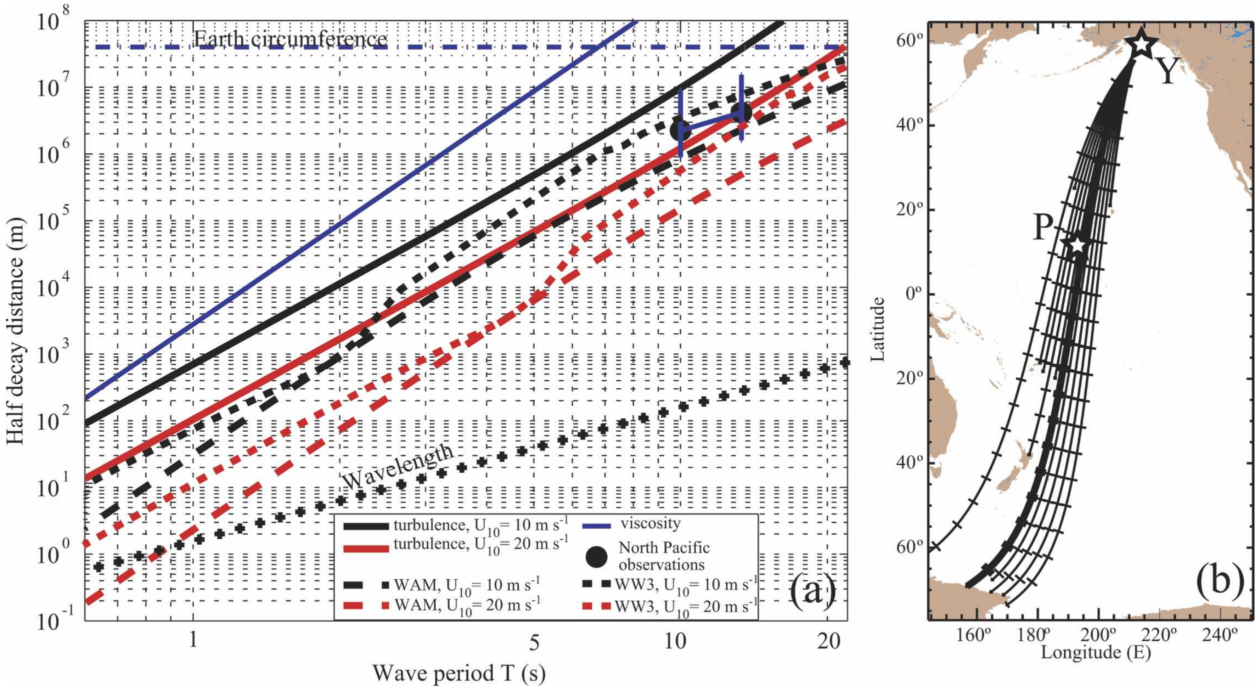

FIG. 2. (a) Spatial scales of wave decay. The present theory for wave attenuation resulting from turbulence for waves of various periods propagating downwind is compared with decay rates observed by Snodgrass et al. (1966) between the Palmyra and Yakutat stations [see (b) for locations], with the dissipation of wave energy attributed to turbulence and whitecapping in the wave model Wavewatch III (WW3: Tolman and Chalikov 1996) or whitecapping only in WAM cycle 4 (Komen et al. 1994), and with viscous dissipation. The following parameter values were used: ρ_a = 1.29 kg m^{-1}, ρ_w = 1026 kg m^{-1}, and $\nu = 3 \times 10^{-6}$ m s^{-2}. Wave model values were obtained by using the linear wave dissipation term after integrating a uniform ocean model over 2 days. The wavelengths (crosses) corresponding to each period, and the earth's circumference are indicated for reference. (b) Example ray trajectories for waves with a period of 12 s arriving in Yakutat, AK, and passing near the island of Palmyra. Transverse tick marks indicate the distance corresponding to one day of propagation. The thick line is the great circle path studied by Snodgrass et al. (1966).

ferent regimes for swells generated at midlatitudes and propagating east in the direction of the dominant winds, versus swells crossing the equator, which may contribute to the difference noted by Snodgrass et al. (1966) between attenuation coefficients over the entire New Zealand–Alaska distance as compared with only the Palmyra–Alaska distance (Fig. 2).

The importance of wave–turbulence interaction can be estimated by calculating the time and length scales of wave decay. We define

$$T_{1/2} = g\frac{2k + 1/\delta}{2k}\frac{\rho_w}{2\rho_a \sigma k u_*^2}\ln 2, \qquad (16)$$

the time it takes for waves to lose one-half of their energy, which would correspond to a 29% reduction in wave height for monochromatic waves. The corresponding propagation distance is $X_{1/2} = C_g/T_{1/2}$, with C_g being the group speed. Here $T_{1/2}$ is about 21 days for an 11-s-period swell with following winds and $u_* = 0.35$ m s^{-1}. Such a value of u_* corresponds to a wind speed of $U_{10} = 10$ m s^{-1}. As the wind speed increases, $T_{1/2}$ and $X_{1/2}$ decrease dramatically (Fig. 2).

For wave periods larger than 6 s, Fig. 2 shows that the present theory gives decay values that are of the same order of magnitude as the dissipation used in the wave-forecasting model Wavewatch III (Tolman and Chalikov 1996), but one order of magnitude smaller than in the cycle 4 of the Wave Model (WAM; Komen et al. 1994). Although the latter model may slightly overestimate dissipation (Rogers et al. 2003) the former probably underestimates the dissipation of waves around the spectral peak because it obtains the same net wave growth while underestimating wind wave generation (see, e.g., Janssen 2004). In addition, the wind-induced attenuation of swells is expected to be several times larger than the effect of turbulence for the lowest-frequency swells (Kudryavtsev and Makin 2004; Ardhuin and Jenkins 2005) so that the swell dissipation in Wavewatch III is probably unrealistically small.

The observations and analysis by Snodgrass et al. (1966) show that waves with periods larger than 15 s were hardly attenuated and waves with periods between 10 and 15 s were not attenuated between New Zealand and the equator but were significantly attenuated from there to Alaska, with a magnitude that is consistent with (15) for following winds of 10–20 m s^{-1} (Fig. 2). Such large mean wind values are highly unlikely over the entire North Pacific Ocean, supporting the idea that other processes also contributed to the wave decay. The dataset of Snodgrass et al. (1966) was insufficient to obtain any significant correlations with the wind direction. New measurements of swell decay using modern technology are needed for further verification.

4. Discussion

Wave shear production of TKE likely accounts for a significant fraction of the energy losses of the wave field, at least for long periods, in addition to the direct production of TKE by breaking waves and attenuation of swell by opposing winds. When applied to the entire wave field, (15) gives a total wave dissipation proportional to the surface Stokes drift $U_s(0)$. Because $U_s(0)$ is the third moment of the surface elevation variance spectrum (the wave spectrum), it is sensitive to the high-frequency waves. As proposed by Kudryavtsev et al. (1999), $U_s(0)$ may be obtained by using a properly defined wave spectrum that matches observations of wave energy (the zeroth moment) and mean square surface slope (the fourth moment). This spectrum gives $U_s(0) = 0.013U_{10}$ for unlimited fetch, which yields a production of TKE of $10u_{*w}^3$, with u_{*w} the water-side friction velocity. This value is about a factor of 10 smaller than the production of TKE usually attributed to wave breaking in active wind wave-generation conditions (e.g., Craig and Banner 1994; Mellor and Blumberg 2004). In addition, this dissipation of wave energy is essentially due to short-period waves that loose most of their energy by breaking (e.g., Melville 1996). Because short waves are more likely to be affected by small-scale currents and correlated with turbulent structures, such as Langmuir circulations, the uniform flux assumption used to derive (13) may be questioned. Yet, there is still no evidence of larger short-wave amplitudes over LC jets (Smith 1980). A proper description of turbulence variations on the scale of the wavelength, including the modulation of short wave breaking by long waves (e.g., Longuet-Higgins 1987), may give a better representation of TKE production.

The present form of (15), already given by McWilliams et al. (1997), should be useful not only for the modeling of upper-ocean mixing, as discussed by Kantha and Clayson (2004), but also for the forecasting of ocean waves and surface drift.

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