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On surface drift currents in the ocean

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A new model of surface drift currents is constructed using the full nonlinear equations of motion. This model includes the balance between Coriolis forces due to the mean and wave-induced motions and the surface wind stresses. The approach used in the analysis is similar to the work by Craik & Leibovich (1976) and Leibovich (1977), but the emphasis is on the mean motion rather than the small-scale time-dependent part of the Langmuir circulation. The final result indicates that surface currents can be generated by both the direct wind stresses, as in the classical Ekman model, and the Stokes drift, derived from the surface wave motion, in an interrelated fashion depending on a wave Ekman number E defined as

$$E = \Omega/\nu_e k_0^2,$$

where Ω is the angular velocity of the earth's rotation, ν_e, the eddy viscosity and k_0, the wavenumber of the surface wave at the spectral peak. When $E \ll 1$, the Langmuir mode dominates. When $E \gg 1$, inertial motion results. The classical Ekman drift current is a special case even under the restriction $E \simeq 1$. On the basis of these results, a new model of the surface-layer movements for future large-scale ocean circulation studies is presented. For this new model both the wind stresses and the sea-state information are crucial inputs.

1. Introduction

The motion of the surface water over the world's oceans is a critical factor in controlling the large-scale transport processes of mass, momentum and energy. It is also the key to solving the global air–sea interaction problem. Past treatment of the surface-water motion has not been totally successful. The solutions obtained reflect the personal preferences of the investigator. The solutions range from the classical Ekman (1905) flow, where currents are generated by the balance between the wind stress and the Coriolis force under a rigid flat surface, to that of Bye (1967) and Kenyon (1970), where currents are attributed to pure Stokes drift from a local wave field subjected to no surface stress. These approaches were not representative of the complete physics. This contention is supported by numerous field observations such as those of Ichiye (1964, 1967), Katz, Gerard & Costin (1965) and Hunkins (1966). The results of Ichiye and Katz *et al.* clearly indicate the existence of an Ekman-type spiral but the shape has a strong dependence on the local sea state. Hunkins' current observations, made under an ice sheet, where no sea-state influence exists, match the expected Ekman spiral to a remarkable degree. Field observations would seem to indicate that any model of surface drift currents devoid of Coriolis and frictional forces must be, in part, fallacious and any model of the surface drift current neglecting the contribution of wave motion would be, at best, incomplete.

The inclusion of wave motion has been attempted by Korvin-Kroukovsky (1972) and Ianniello & Garvine (1975) using uncoupled models in which the total drift currents are calculated as the sum of the individual Stokes and Ekman components. Their conclusions are that the wave-induced drift is dominant. Unfortunately, these uncoupled models fail to consider the influences of the Coriolis forces generated by the Stokes drift and the interaction between the Ekman drift and the wave motions. Consequently, the results are questionable.

A considerable advance in the problem of the inclusion of wave motion has been made by Craik & Leibovich (1976) and Leibovich (1977). Although the Coriolis force was neglected, they achieved a major breakthrough by inclusion of the wave motion through a rigorous procedure where the Stokes drift is represented in an Eulerian framework. In the present analysis, the formulation is based on that of Leibovich (1977) with the essential modification of adding the Coriolis term so that the Ekman drift can be included. This modification does not invalidate the results of Craik & Leibovich (1976) and Leibovich (1977), which are primarily for the small-scale motions, but rather extends their work to cover large-scale mean motion. The results of this model indicate that the sea state could have strong influence on surface drift currents in the ocean. Furthermore, a classification scheme is proposed to explain the highly variable conditions of the surface drift currents.

2. Analysis

Within the oceanic surface layer, the predominant motions of the water are due to gravity waves. Since these waves can be successfully approximated by an irrotational motion, the total velocity field $\mathbf{q}'$ can be expressed as the sum of the velocity $\mathscr{U}'$ associated with the linearized irrotational wave motion and the higher-order velocity perturbation $\mathbf{v}'$ caused by waves and wind, i.e.

$$\mathbf{q}' = \epsilon\mathscr{U}' + \epsilon^2\mathbf{v}', \tag{1}$$

in which ϵ is the perturbation parameter, assumed to be of the order of the surface wave slope. Using this parameter, the higher-order velocity $\mathbf{v}'$ can be further expanded as

$$\mathbf{v}' = \mathbf{v}_0' + \epsilon\mathbf{v}_1' + \epsilon^2\mathbf{v}_2' + \dots. \tag{2}$$

Assuming an incompressible fluid, we can write the equations of motion as

$$\partial\mathbf{q}'/\partial t' + \mathbf{q}'.\nabla\mathbf{q}' + 2\mathbf{\Omega}\times\mathbf{q}' = -\rho^{-1}\nabla p' + \nu_e\nabla^2\mathbf{q}', \tag{3}$$

$$\nabla.\mathbf{q}' = 0, \tag{4}$$

where ν_e is the eddy viscosity and Ω is the angular velocity of the earth's rotation. The vorticity equation can then be formed by taking the curl of (3) to eliminate the pressure term. This results in

$$\partial\boldsymbol{\varpi}'/\partial t' + \nabla\times(\boldsymbol{\varpi}'\times\mathbf{q}') + 2\nabla\times(\mathbf{\Omega}\times\mathbf{q}') = \nu_e\nabla^2\boldsymbol{\varpi}', \tag{5}$$

where

$$\boldsymbol{\varpi}' = \nabla\times\mathbf{q}' = \epsilon^2\boldsymbol{\varpi}_0' + \epsilon^3\boldsymbol{\varpi}_1' + \epsilon^4\boldsymbol{\varpi}_2' + \dots. \tag{6}$$

In order to scale the various terms in (5) properly a length scale $1/k_0$ and a time scale $1/\sigma_0$ are introduced, where k_0 and σ_0 are the wavenumber and frequency of the surface gravity waves at the spectral peak, respectively. With this choice of scales, the vorticity equation (5) becomes

$$\frac{\partial \boldsymbol{\varpi}}{\partial t} + \nabla \times (\boldsymbol{\varpi} \times \mathbf{q}) + \frac{2\Omega}{\sigma_0} \nabla \times (\mathbf{e} \times \mathbf{q}) = \frac{\nu_e k_0^2}{\sigma_0} \nabla^2 \boldsymbol{\varpi}, \tag{7}$$

where the primes are dropped to indicate non-dimensional quantities and Ω is replaced by $\Omega\mathbf{e}$, with $\mathbf{e}$ a unit vector parallel to the axis of the earth's rotation.

Next, the motion will be divided into mean and fluctuating parts as

$$\mathbf{v} = \bar{\mathbf{v}} + \langle \mathbf{v} \rangle, \quad \boldsymbol{\varpi} = \bar{\boldsymbol{\varpi}} + \langle \boldsymbol{\varpi} \rangle, \tag{8}$$

the overbars indicating the time-averaged values and the angular brackets the fluctuating parts. The perturbed equations of motion to order ϵ^2, ϵ^3 and ϵ^4, respectively, can then be written as

$$\partial \boldsymbol{\varpi}_0 / \partial t = 0, \tag{9}$$

$$\frac{\partial \boldsymbol{\varpi}_1}{\partial t} + \nabla \times (\boldsymbol{\varpi}_0 \times \boldsymbol{\mathscr{U}}) - \frac{2\Omega}{\sigma_0 \epsilon^2} (\mathbf{e} \,.\, \nabla) \boldsymbol{\mathscr{U}} = 0, \tag{10}$$

and

$$\frac{\partial \boldsymbol{\varpi}_2}{\partial t} + \nabla \times (\boldsymbol{\varpi}_0 \times \mathbf{v}_0) + \nabla \times (\boldsymbol{\varpi}_1 \times \boldsymbol{\mathscr{U}}) - \frac{2\Omega}{\sigma_0 \epsilon^2} (\mathbf{e} \,.\, \nabla) \mathbf{v}_0 = \frac{\nu_e k_0^2}{\sigma_0 \epsilon^2} \nabla^2 \boldsymbol{\varpi}_0. \tag{11}$$

Equations (9), (10) and (11) are similar to those of Leibovich (1977) except for the Coriolis terms. From (9), it can be shown that

$$\langle \boldsymbol{\varpi}_0 \rangle = 0, \tag{12}$$

because the fluctuating part of $\boldsymbol{\varpi}_0$ is induced by the periodic motions. Note, however, that (12) does not imply that $\boldsymbol{\varpi}_0 = 0$.

It then follows from (10) that

$$\begin{aligned} \langle \boldsymbol{\varpi}_1 \rangle &= \nabla \times \left(\int^t \boldsymbol{\mathscr{U}} d\tau \times \bar{\boldsymbol{\varpi}}_0 \right) + \frac{2\Omega}{\sigma_0 \epsilon^2} (\mathbf{e} \,.\, \nabla) \int^t \boldsymbol{\mathscr{U}} d\tau \\ &= (\bar{\boldsymbol{\varpi}}_0 \,.\, \nabla) \int^t \boldsymbol{\mathscr{U}} d\tau - \left(\int^t \boldsymbol{\mathscr{U}} d\tau \,.\, \nabla \right) \bar{\boldsymbol{\varpi}}_0 + \frac{2\Omega}{\sigma_0 \epsilon^2} (\mathbf{e} \,.\, \nabla) \int^t \boldsymbol{\mathscr{U}} d\tau. \end{aligned} \tag{13}$$

Next, taking the mean of (11) yields

$$\nabla \times (\bar{\boldsymbol{\varpi}}_0 \times \bar{\mathbf{v}}_0) + \nabla \times \overline{(\boldsymbol{\varpi}_1 \times \boldsymbol{\mathscr{U}})} - \frac{2\Omega}{\sigma_0 \epsilon^2} (\mathbf{e} \,.\, \nabla) \bar{\mathbf{v}}_0 = \frac{\nu_e k_0^2}{\sigma_0 \epsilon^2} \nabla^2 \bar{\boldsymbol{\varpi}}_0, \tag{14}$$

with

$$\overline{\nabla \times (\boldsymbol{\varpi}_1 \times \boldsymbol{\mathscr{U}})} = \overline{(\boldsymbol{\mathscr{U}} \,.\, \nabla) \langle \boldsymbol{\varpi}_1 \rangle} - \overline{(\langle \boldsymbol{\varpi} \rangle \,.\, \nabla) \boldsymbol{\mathscr{U}}}. \tag{15}$$

Then, combining (13) and (15), and using tensor notation for convenience, the result is

$$\begin{aligned} \{\nabla \times \overline{(\boldsymbol{\varpi}_1 \times \boldsymbol{\mathscr{U}})}\}_i &= \overline{\mathscr{U}_k \left\{ \varpi_{0j} \int^t \mathscr{U}_{i,j} d\tau \right\}_{,k}} - \overline{\mathscr{U}_k \left\{ \varpi_{0i,j} \int^t \mathscr{U}_j d\tau \right\}_{,k}} \\ &+ \frac{2\Omega}{\sigma_0 \epsilon^2} \overline{\mathscr{U}_k \left\{ e_j \int^t \mathscr{U}_{i,j} d\tau \right\}_{,k}} - \overline{\varpi_{0k} \mathscr{U}_{i,j} \int^t \mathscr{U}_{j,k} d\tau} + \overline{\varpi_{0k,j} \mathscr{U}_{i,k} \int^t \mathscr{U}_j d\tau} \\ &- \frac{2\Omega}{\sigma_0 \epsilon^2} e_j \overline{\int^t \mathscr{U}_{k,j} d\tau \, \mathscr{U}_{i,k}}. \end{aligned} \tag{16}$$

Observe that the Stokes drift $\mathscr{U}_s$ can be written as

$$\mathscr{U}_{si} = \overline{\mathscr{U}_{i,j}\int^t \mathscr{U}_j\,d\tau}. \tag{17}$$

Following the scheme used by Craik & Leibovich (1976), we can write the Coriolis terms in (16) as

$$\begin{aligned}
&\frac{2\Omega}{\sigma_0\epsilon^2}e_j\left\{\overline{\mathscr{U}_k\int^t \mathscr{U}_{i,jk}\,d\tau} - \overline{\mathscr{U}_{i,k}\int^t \mathscr{U}_{k,j}\,d\tau}\right\}\\
&= \frac{2\Omega}{\sigma_0\epsilon^2}e_j\left\{-\mathscr{U}_{si,j} + \overline{\mathscr{U}_{i,kj}\int^t \mathscr{U}_k\,d\tau} + \overline{\mathscr{U}_k\int^t \mathscr{U}_{i,jk}\,d\tau}\right\}\\
&= -\frac{2\Omega}{\sigma_0\epsilon^2}e_j\mathscr{U}_{si,j} + \frac{2\Omega}{\sigma_0\epsilon^2}e_j\frac{\partial}{\partial t}\overline{\left\{\int^t \mathscr{U}_{i,kj}\,d\tau\int^t \mathscr{U}_k\,d\tau\right\}}\\
&= -\frac{2\Omega}{\sigma_0\epsilon^2}e_j\mathscr{U}_{si,j}.
\end{aligned} \tag{18}$$

Then by combining (14), (16) and (18), the following is obtained:

$$-\frac{\nu_e k_0^2}{\sigma_0\epsilon^2}\nabla^2\bar{\boldsymbol{\varpi}}_0 = (\bar{\boldsymbol{\varpi}}_0\,.\,\nabla)(\bar{\mathbf{v}}_0+\mathscr{U}_s) - (\bar{\mathbf{v}}_0+\mathscr{U}_s)\,.\,\nabla\bar{\boldsymbol{\varpi}}_0 + \frac{2\Omega}{\sigma_0\epsilon^2}(\mathbf{e}\,.\,\nabla)(\bar{\mathbf{v}}_0+\mathscr{U}_s). \tag{19}$$

This is the same expression as equation (14) in Craik & Leibovich (1976) with the addition of the extra term representing Coriolis forces. Equation (19) is the generalized Ekman equation with wave motion included.

3. Specific results

Having derived the generalized Ekman equation, we can seek an Ekman-type solution by assuming that all the mean motions are functions of z alone; then the relations for the velocity components can be written as

$$\frac{\partial^3\bar{v}_0}{\partial z^3} = 2\tilde{E}\,\frac{\partial(\bar{U}_0+\mathscr{U}_s)}{\partial z}, \quad \frac{\partial^3\bar{U}_0}{\partial z^3} = -2\tilde{E}\,\frac{\partial(\bar{v}_0+\mathscr{V}_s)}{\partial z}, \tag{20}$$

where $\tilde{E} = f/\nu_e k_0^2$ is an Ekman-type number, with $f = \mathbf{\Omega}\,.\,\mathbf{e}_3$, the local component of the earth's rotation, and $\mathbf{e}_3$ the unit vector in the local vertical direction. The significance of the Ekman-type number $\tilde{E}$ (or more generally, $E \equiv \Omega/\nu_e k_0^2$) will be discussed in detail later.

For a random gravity wave field, the Stokes drift can be expressed, as in Huang (1971), as

$$\mathscr{U}_s' = \int_{\mathbf{k}}\int_n 2n\mathbf{k}\chi(\mathbf{k},n)\exp(2|\mathbf{k}|z)\,d\mathbf{k}\,dn, \tag{21}$$

where $\mathbf{k}$ is the wavenumber vector, n is the frequency and $\chi(\mathbf{k}, n)$ is the directional wave energy spectrum. If we define the current at the surface as $\tilde{\mathbf{Q}}_0'$, then the solution of (20) expressed in dimensional form will be

$$\begin{aligned}
\bar{\mathbf{Q}}_0' &= \tilde{\mathbf{Q}}_0'\exp\{(f/\nu_e)^{\frac{1}{2}}(1+i)z'\}\\
&\quad + i\int_{\mathbf{k}}\int_n \frac{2n\mathbf{k}}{(2\nu_e/f)|\mathbf{k}|^2-i}\left\{\exp 2|\mathbf{k}|z' - \exp\left(\frac{f}{\nu_e}\right)^{\frac{1}{2}}(1+i)z'\right\}\chi(\mathbf{k},n)\,d\mathbf{k}\,dn,
\end{aligned} \tag{22}$$

where the relationship $\bar{\mathbf{Q}}_0' = \bar{U}_0' + i\bar{V}_0'$ holds.

If we consider the specific case when the wind is blowing in the $+y$ direction, then the surface boundary conditions on (20) can be specified exactly as in Ekman's original paper of 1905, i.e.

$$\partial U_0/\partial z = 0, \quad \partial V_0/\partial z = S/\sigma_0 \rho \nu_e \quad \text{at} \quad z = 0, \tag{23}$$

where S is the surface stress and ρ is the density of the sea water. The solutions of (20) in component form and dimensional variables can easily be shown to be

$$\bar{U}_0' = C_1 \exp(az') \cos(az' + C_2) - \int_{\mathbf{k}} \int_n \frac{2n[k_x + 2k_y |\mathbf{k}|^2/a^2]}{(2|\mathbf{k}|^2/a^2)^2 + 1} \exp(2|\mathbf{k}|z') \chi(\mathbf{k}, n) d\mathbf{k}\, dn, \tag{24}$$

$$\bar{V}_0' = C_1 \exp(az') \sin(az' + C_2) + \int_{\mathbf{k}} \int_n \frac{2n[2k_x |\mathbf{k}|^2/a^2 - k_y]}{(2|\mathbf{k}|^2/a^2)^2 + 1} \exp(2|\mathbf{k}|z') \chi(\mathbf{k}, n)\, d\mathbf{k}\, dn, \tag{25}$$

where $a^2 = f/\nu_e$, $\mathbf{k} = k_x + ik_y$ and C_1 and C_2 are given by

$$C_2 = \arctan \frac{\dfrac{S}{\nu_e \rho a} - \displaystyle\int_{\mathbf{k}} \int_n \frac{(k_x - k_y) + 2(k_x + k_y)|\mathbf{k}|^2/a^2}{(2|\mathbf{k}|^2/a^2)^2 + 1} \frac{4n|\mathbf{k}|}{a} \chi(\mathbf{k}, n)\, d\mathbf{k}\, dn}{\dfrac{S}{\nu_e \rho a} - \displaystyle\int_{\mathbf{k}} \int_n \frac{(k_x + k_y) - 2(k_x - k_y)|\mathbf{k}|^2/a^2}{(2|\mathbf{k}|^2/a^2)^2 + 1} \frac{4n|\mathbf{k}|}{a} \chi(\mathbf{k}, n)\, d\mathbf{k}\, dn}, \tag{26}$$

and

$$C_1 = \frac{1}{\cos C_2 + \sin C_2} \left\{ \frac{S}{\nu_e \rho a} - \int_{\mathbf{k}} \int_n \frac{[2k_x |\mathbf{k}|^2/a^2 - k_y]}{(2|\mathbf{k}|^2/a^2)^2 + 1} \frac{4n|\mathbf{k}|}{a} \chi(\mathbf{k}, n)\, d\mathbf{k}\, dn \right\}. \tag{27}$$

Some special cases will be considered. The first case is a sea surface where surface waves are absent. The directional wave energy spectrum $\chi(\mathbf{k}, n)$ is identically zero:

$$\chi(\mathbf{k}, n) \equiv 0. \tag{28}$$

Substitution of (28) into (26) and (27) yields

$$C_2 = \tfrac{1}{4}\pi, \quad C_1 = S/2^{\frac{1}{2}} \nu_e \rho a. \tag{29}$$

These constants of integration are identical to those first presented by Ekman (1905).

The second case is a sea surface with the local waves either parallel or symmetric with respect to the local wind, which is oriented along the y axis. Then, as previously shown, $C_2 = \frac{1}{4}\pi$ but now the surface currents can move in any direction between 0 and $\frac{1}{2}\pi$ according to the relative magnitude of U_0 and V_0. This may provide an explanation for the directional variation of surface drift currents observed in field data.

4. Discussion

An interesting feature of the generalized Ekman equation (19) is that, although the Stokes drift contributes to the Coriolis force it does not appear in the viscous term. This is reasonable because the Stokes drift is a consequence of nonlinear effects of inviscid waves. The inclusion of the Stokes drift in the generalized Ekman equation, however, provides the necessary coupling between the sea state and the Ekman flow. The result of the present analysis indicates that, for a realistic estimation of the water mass movement at the surface layer, one needs not only information on the wind stresses but also information on the sea state, in the form of the directional spectrum.

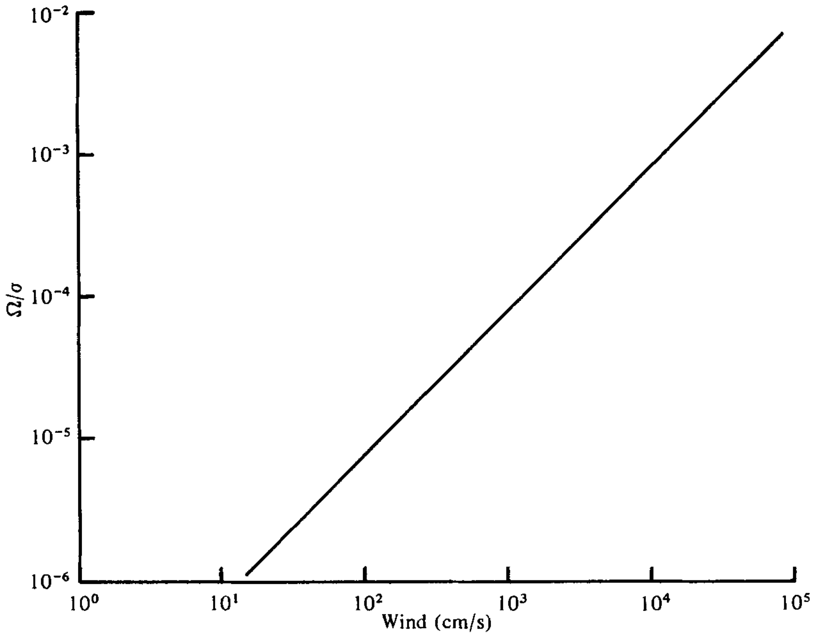

FIGURE 1. Variation of Ω/σ with wind velocity.

The inclusion of the sea state may yield new insight into the study of large-scale air–sea interaction and a more complete understanding of ocean circulation.

The expression given in (22) links the sea-state condition with the fundamental Ekman circulation. This was achieved under a rather restrictive Ekman-type assumption, i.e. that all the motions are horizontal and all the functions depend on z alone. Nevertheless, the rigorous development presented in this analysis enables a discussion of the more general conditions of the surface drift currents under the influence of both the wave motion and the wind stress.

For more detailed physical discussions, examine the generalized Ekman equation (19). In the derivation of this equation the scaling of the terms $\Omega/\sigma_0\epsilon^2$ and $\nu_e k_0^2/\sigma_0\epsilon^2$ must be comparable with the rest of the terms, or the whole analysis would be wrong. Assume that the waves are all wind generated, then the dominant wave frequency σ_0 can be related to the wind velocity W (see, for example, Phillips 1966) by

$$\sigma_0 = g/W.$$

This gives

$$\Omega/\sigma_0\epsilon^2 = \Omega W/g\epsilon^2. \tag{30}$$

Figure 1 presents Ω/σ_0 or $\Omega W/g$ as a function of the wind speed. For a typical wind of 10 m/s, Ω/σ_0 is the order of 10^{-4}. This requires an ϵ of the order of 10^{-2} to make the $\Omega/\sigma_0\epsilon^2$ terms comparable to the other terms in the generalized Ekman equation. The stability limitation on the gravity waves allows wave slopes up to the order of 10^{-1}, but this value applies only to the higher wavenumbers, where the individual waves are actively breaking. The main energy-containing components over most of the open ocean are far more gentle than the breaking waves. On the basis of the most recent JONSWAP data reported by Hasselmann *et al.* (1973, 1976), the mean slope of the waves, defined by $\overline{\zeta^2}k^2$, with $k = g/W^2$, is around 10^{-3}–10^{-5}. Next we have to

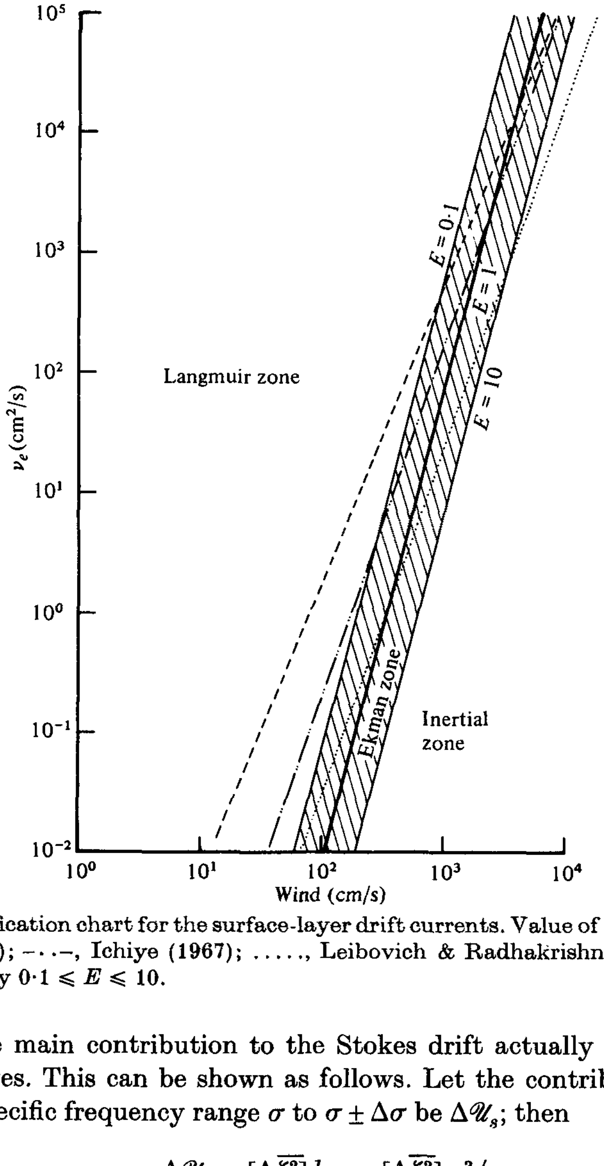

FIGURE 2. Classification chart for the surface-layer drift currents. Value of ν_e from: – – –, Neumann & Pierson (1964); –·· –, Ichiye (1967); ....., Leibovich & Radhakrishnan (1977). The Ekman zone is defined by $0{\cdot}1 \leqslant E \leqslant 10$.

prove that the main contribution to the Stokes drift actually comes from the low frequency waves. This can be shown as follows. Let the contribution to the Stokes drift from a specific frequency range σ to $\sigma \pm \Delta\sigma$ be $\Delta\mathscr{U}_s$; then

$$\Delta\mathscr{U}_s = [\Delta\overline{\zeta^2}]\,k\sigma = [\Delta\overline{\zeta^2}]\,\sigma^3/g, \tag{31}$$

where $[\Delta\overline{\zeta^2}]$ represents the energy in the frequency band σ to $\sigma \pm \Delta\sigma$. The dispersion relationship has been used in (31). By using the equilibrium form of the spectrum proposed by Phillips (1958), $[\Delta\overline{\zeta^2}]$ can be written as

$$[\Delta\overline{\zeta^2}] = \frac{\alpha g^2}{\sigma^5}\,2\Delta\sigma, \tag{32}$$

where the non-dimensional constant α is of the order of 10^{-2}. Now, combining (31) and (32), we get

$$\Delta\mathscr{U}_s = \frac{\alpha g}{\sigma^2}\,2\Delta\sigma.$$

Locality	Layer	ν_e(cm²/s)	Source	Author(s)
(I) *Absolute constants*				
Danish Waters	0–15 m	1·9–3·8	All currents	Jacobsen (1913)
Arctic Ocean		160	Under ice	Brennecke (1921)
Danish Waters		250–1500	All currents	Durst (1924)
Kuroshio	0–200 m	680–7500	All currents	Suda (1926)
Japan Sea	0–200 m	150–1460	All currents	Suda (1926)
North Siberian Shelf	0–60 m	0–1000	Tidal currents	Sverdrup (1926)
North Sea	0–31 m	75–1720	Strong tidal currents	Thorade (1928)
Tropical Atlantic Ocean	0–50 m	320	Temperature fluctuation	Defant (1932)
North Siberian Shelf	0–60 m	10–400	Tidal current	Fjeldstad (1936)
Atlantic Ocean 50° S–10°N	0–200 m	7–50	Wind currents	Wüst (1955)
Japan Sea	0–10 m	100	Tidal currents	Nan'niti (1964)
Arctic Ocean	0–100 m	23·8	Ice drift	Hunkins (1966)
North Atlantic	0–4 m	146	Temperature fluctuation	Bowden, Howe & Tait (1970)
	4–8 m	96		
	8–12 m	47		
Open Ocean	0–10 m	150–225	Surface	Assaf, Gerard & Gordon (1971)
Lake Huron	30 m	65–160	Wind current	Csanady (1972)
Tropical Atlantic Ocean	0–12 m	420 ± 84	Temperature fluctuation	Hoeber (1972)
Tropical Atlantic Ocean	0–10 m	62	Temperature fluctuation	Ostapoff & Worthem (1974)
	10–20 m	68		
	20–30 m	85		
Tropical Atlantic Ocean	0–12 m	480	Temperature fluctuation	Prümm (1974)
	20–50 m	265		
(II) *Relative constants*				
All Oceans	Surface	$1{\cdot}02\,W^3$ for $W < 6$ m/s $4{\cdot}30\,W^2$ for $W > 6$ m/s	Thickness of upper homogeneous layer	Ekman (1905), Thorade (1914)
Theoretical	All	$l^{4/3}\epsilon^{1/3}$		Kolmogorov (1941)
All Oceans	Surface	$5 \times 10^5 ca$	Wave motion	Bowden (1950)
Atlantic Ocean	Surface	$0{\cdot}1825 \times 10^{-4} W^{5/2}$		Neumann & Pierson (1964)
Liverpool Bay		$2{\cdot}5 \times 10^{-3} Uh$	Tidal currents	Bowden (1967)
All Oceans	Surface	$\frac{1}{f}\left(\frac{\rho_w}{\rho_0}\frac{c_{10}}{\kappa}\right)^2 W_{10}^2$		Kullenberg (1976)
Wangara Data	All	$\frac{1}{20}u_* h$	Shallow current	Csanady (1976)
	Surface	$u_*^2/200f$	Deep ocean	
Ocean surface	Surface	$(2\text{–}6) \times 10^{-5} W^3/g$	Surface drift	Leibovich & Radhakirshnan (1977)

		(III) *Implicit function*		
Baltic Sea	All	$\nu_0(1+\beta Ri)^{-\frac{1}{2}}$	All currents	Munk & Anderson (1948)
Theoretical	All	$\kappa u_* L R_f$	All	Monin & Oboukhov (1954)
All oceans	All	$\nu_0 \exp(-mRi)$	All	Mamayev (1958)
Laboratory	All	$R_i^{-\frac{3}{2}}$	Wind data	Webster (1964)
All oceans	All	$8{\cdot}9 \times 10^{-8}\, W^2 N^{-2} \lvert dU/dz \rvert$	All	Kullenberg (1971)
All oceans	All	$(1-r\theta')^{-1}$	Theory	Kuftarkov & Fel'zenbaum (1976)
		(IV) *Explicit function*		
North Siberian Shelf	0–22 m	$385\left(\frac{z+0{\cdot}1}{22{\cdot}1}\right)$	Wind currents	Fjeldstad (1929)
Theoretical	All	$\kappa u_*(z+z_0)$	All	Von Kármán (1930), Thomas (1975)
Germany	Surface	$0{\cdot}02(z+z_0)$	Wind data	Mildner (1932)
	Surface	$(f/3\sqrt{2})(h-z)^2$	Wind data	Rossby & Montgomery (1935)
Theoretical	Surface	$\kappa(z+S_w\epsilon)(\rho_a/\rho)\, rW_0$	Wind current	Rossby & Montgomery (1935)
Theoretical	All	$\kappa(z+z_0)\, W_a/\ln[(z_a+z_0)/z_0]$	Air flow	Rossby & Montgomery (1936)
Theoretical	Surface	$\frac{\pi k^2}{18}\frac{H^2}{T}\exp(-2kz)[1-\pi^2\delta^2\exp(2kz_1)]^3$	Waves	Dobrokinoskii (1947)
Theoretical	All	$\kappa u_* z/(1+\alpha z/L)$	Air	Monin & Oboukhov (1954)
Theoretical	Surface	$(0{\cdot}005/4\pi)\, gTH \exp(-kz)$	Waves	Bashkirov (1959)
Theoretical	Surface	$a^3 k\sigma \exp(-kz)$	Waves	Kitaigorodskii (1960)
Theoretical	All	$\kappa u_* z/L[(1-\exp(-z/L)]$	Air	Swinbank (1964)
Theoretical	All	$\kappa u_* z\lambda(R_f)(1-R_f)^{\frac{1}{4}}$	Air	Monin & Yaglom (1971)
Theoretical	All	$\kappa u_* z/(1-rz/L)^{\frac{1}{4}}$	Air	Yamamoto & Shimasuki (1966), Businger (1966)
New York Bight	50 m	$0{\cdot}028(H^2/T)\exp(-2kz)$	Dye motion	Ichiye (1967)
Theoretical	Surface	$\tilde{\kappa} u_* L_*$ $\tilde{\kappa} = \frac{(96)^{\frac{1}{3}} B^3}{\delta^{\frac{4}{3}}} F^{\frac{1}{2}}(\tilde{z})\exp(-2^{\frac{1}{2}} B\tilde{z}^{\frac{1}{2}})$	Waves	Kitaigorodskii & Miropolskey (1968)
Theoretical	All	$\kappa u_* z/(1-16Ri)^{\frac{1}{4}}$	All	Miyake *et al.* (1970)
Theoretical	Surface	$\nu_0 \exp(kz)$	Surface	Witten & Thomas (1976)
Theoretical	Surface	$E_T^{\frac{1}{2}}\, l$	Wave motion	Benilov (1973)

TABLE 1. Values of the vertical eddy viscosity ν_e.

Hence the contribution to the Stokes drift from a specific frequency band is inversely proportional to the square of the frequency. Therefore the low frequency waves are more important in the drift-current generation. For these low frequency waves, it is not unreasonable to use $\epsilon^2 = O(10^{-4})$. Consequently, the assumption $\Omega/\sigma_0 \epsilon^2 = O(1)$ is well within reasonable limits.

Having established the order of Ω/σ_0, we can discuss the magnitude of the viscous term by forming the ratio of the two terms. This results in a wave-related, Ekman-number-like parameter E defined as

$$E = \Omega/\nu_e k_0^2.$$

For a wind wave field, k_0 can be related to the wind field as $k_0 = g/W^2$, then

$$E = \Omega W^4/\nu_e g.$$

The value of E is plotted in figure 2 as a contour map in ν_e, W space. For a typical wind speed of 10 m/s, $E = 1$ requires a value of ν_e of 75 cm^2/s; but for $E = O(1)$, a range of ν_e of 10–1000 cm^2/s is satisfactory. These values are all well within the range of commonly adopted ν_e values. Thus, under most natural conditions, the surface flow will have an E of order one, i.e. the viscous term and the Coriolis term are of the same order. Since the wind conditions and the relationship between the wind and the wavenumber of the energy-containing waves are all well defined, the detailed discussion of the surface drift will hinge on the value of the eddy viscosity.

The determination or the parameterization of the eddy viscosity ν_e in terms of observable physical quantities is one of the most difficult problems in physical oceanographic studies. A list of the commonly used ν_e values is given in table 1, where the ν_e values and/or the parameterized forms of ν_e are grouped by their characteristic properties.

Since the eddy viscosity is no longer a physical property of the fluid but rather a dynamic property of the specific flow, the constant values in group I can not be very meaningful or representative. The second group shows the values changing with environmental conditions, but remaining relatively constant throughout the flow field. Arguments against using these values are obvious from the fact that the upper ocean layer has inhomogeneous vertical temperature, salinity and turbulent intensity stratification. The last two groups list the values of the eddy viscosity as implicit or explicit functions of spatial variables and environmental conditions. These functions are the most reasonable expressions, but short of definitive proof, their application will add unnecessary complications to the problem. The principle of a vertically variable eddy viscosity is essential for any realistic surface drift current model. This will be discussed in the following sections.

For the sake of simplicity, some of the values in the second group will be used as examples. Typical of the expressions for the eddy viscosity is the empirical formula given by Neumann & Pierson (1964):

$$\nu_e = 0{\cdot}1825 \times 10^{-4} W^{\frac{5}{2}},$$

where W is the wind speed in cm/s and ν_e is in cm^2/s. Since it is dimensionally incorrect, this expression cannot be very general physically, yet it is widely used by oceanographers. A second expression, from Leibovich & Radhakrishnan (1977), is

$$\nu_e = 2{\cdot}84 \times 10^{-5}\, W^3/g.$$

Both of these expressions relate ν_e to the surface wind speed directly. The second, being consistent dimensionally, is more meaningful dynamically. The third expression relating ν_e to the sea state is a modified version of the expression proposed by Ichiye (1967):

$$\nu_e = 0{\cdot}028\, H^2/T,$$

where H is the wave height and T is the wave period. This viscosity can also be related to the wind speed.

Figure 2 shows the values of ν_e predicted by the three expressions. It is clear that under natural wind conditions, ranging from a few m/s to a few tens of m/s, the surface drift current should be controlled equally by Coriolis and frictional forces.

Not only is the value of ν_e important in determining the characteristics of drift currents, but the vertical variation of ν_e is also critical in constructing the detailed model of the surface drift current structure. Unfortunately, detailed knowledge of the vertical variation of ν_e is still lacking. The most commonly accepted form for ν_e is a linearly increasing function of depth. This is based on an inverted atmospheric boundary-layer model where the velocity profile is given by the logarithmic function

$$W(z) = \frac{W_*}{\kappa} \ln \frac{z}{z_0}, \tag{33}$$

in which W_* is the frictional velocity κ is the von Kármán constant and z_0, is a roughness parameter. With the velocity profile given as (33), and the constant-stress assumption, upon which (33) is based, it is easy to show that

$$\nu_e = \kappa W_* z. \tag{34}$$

Admittedly, the study of atmospheric mixing is far more advanced than its oceanic counterpart, yet this indiscrimate borrowing of the atmospheric result is hardly justifiable. In the atmosphere the ground acts as a barrier to the flux of momentum. Consequently the intensity and the scale of turbulence decrease as the height from the surface decreases. On the other hand, the ocean surface, which is constantly acted on by the wind stresses, breaking waves, etc. is a source of turbulence. Thus the turbulent intensity in the ocean surface layer should be stronger and the mixing more thorough, as manifested by the existence of the homogeneous upper layer above the thermocline. If one accepts the existence of the mixed layer at the top of the upper ocean as evidence of strong mixing, then the eddy viscosity should be represented by a decreasing function of depth rather than an increasing function.

Under these assumptions, the drift current in the top layer can be classified according to the dominating mechanism of the motion, as follows. For a given wind condition, the turbulent intensity in the top surface layer, being strongly influenced by the active breaking of waves, is high, so the value of ν_e will also be high. Then E will be small. Consequently, the motion is more likely to be in the Langmuir zone. The thickness of the layer dominated by the Langmuir mechanism is strictly determined by the turbulent intensity of the high frequency breaking waves convected by the orbital velocities of the large waves. Therefore the dominance of the Langmuir mechanism cannot exceed a few amplitudes of the main energy-containing waves. This, however, does not imply the cessation of all the vertical motion associated with the Langmuir cells. In fact, according to Leibovich (1977), the vertical motion can extent to much greater depths. As the turbulence decays with the depth so does ν_e.

Author(s)	Test conditions	Deflexion angle, θ	Drift current, $\|V_0\|$
Ekman (1905)	Theoretical	45°	$0{\cdot}0127\ W/\sin^{\frac{1}{2}}\phi$
Mohn (1883)			$0{\cdot}0103\ W$
Dinklage (1888)			$0{\cdot}0127\ W$
Nansen (1902)	Ice floats	20–40°	$0{\cdot}019\ W$
Forch (1909)	Mediterranean	$44{\cdot}8 \pm 21°$	
Witting (1909)	Finish light ship	$34\text{–}7{\cdot}5°(W\ \mathrm{m/s})^{\frac{1}{2}}$	$0{\cdot}010\ W$
Gallé (1910)	Indian Ocean	$47{\cdot}3 \pm 7{\cdot}3°$	
Thorade (1914)			$2{\cdot}59(W/\sin\phi)^{\frac{1}{2}}$ for $W \leqslant 6\ \mathrm{m/s}$ $1{\cdot}26(W/\sin\phi)^{\frac{1}{2}}$ for $W > 6\ \mathrm{m/s}$
Brennecke (1921)	Ice floats		$0{\cdot}0269\ W$
Durst (1924)			$0{\cdot}0079\ W/\sin^{\frac{1}{2}}\phi$
Sverdrup (1928)			$0{\cdot}0177\ W$
Palmén (1930, 1931)	8 m/s wind	$9{\cdot}1 \pm 4{\cdot}8°$	$0{\cdot}0114\ W$
Smith (1931)	Iceberg drift	40° (deep immersing iceberg) 18–21° (small iceberg)	$0{\cdot}006\ W$ $0{\cdot}012\ W$
Rossby (1932)	Theory	54° 44′	$\frac{\gamma}{\kappa}\left(\frac{2\rho_a}{3\rho_w}\right)^{\frac{1}{2}} W$
Rossby & Montgomery (1935)	Theory of ice drift	$\cot^{-1}\left\{0{\cdot}586\left[2{\cdot}835+\ln\frac{W}{fz_0}\,10^{-6}\right]\right\}$	$1{\cdot}8\times10^{-2}\ W/\sin\phi$
Rossby & Montgomery (1935)	Theory on water	$\tan^{-1}[2/(\sqrt{2}+3z)]$ $z=\frac{\kappa}{\kappa_0}\ln\left(\frac{3\gamma k^2}{k_0 S_w \epsilon f}\left(\frac{\rho}{2\rho_w}\right)^{\frac{1}{2}} W\right)$	$\frac{\gamma}{\kappa}\left\{\frac{\rho}{\rho_w}\left(\frac{2}{3}+\frac{2\sqrt{2z}}{3}+z^2\right)^{\frac{1}{2}} W\right.$

Neumann (1939)		$22\text{–}6{\cdot}3^\circ\,[(W-4)\,\mathrm{m/s}]^{\frac{1}{2}}$	
Van Dorn (1953)	Pond	0	$k_1\sqrt{R} : R \ll 10^3$ $0{\cdot}02\,W : R \gg 10^3$
Stommel (1954)	Off Bermuda	$37{\cdot}9 \pm 31{\cdot}5^\circ$	CW or $C_1\,W^{\frac{3}{2}}$
Faller (1964)	Drift cards	$12{\cdot}1 \pm 6{\cdot}5^\circ$	
Tomczak (1964)	Cards	0	$0{\cdot}042\,W$
Sutcliffe, Baylor & Menzel (1963)	Foam, slick, weed	$14 \pm 60^\circ$	
Ichiye (1964)	Theory	$q = i2Tz\beta^{-1}(m_2 n_1 - m_1 n_2)^{-1}\{n_2 H_1^{(1)}(\beta z) - n_1 H_1^{(2)}(\beta z)\}$	
Hunkins (1966)	Ice drift	$34{\cdot}5 \pm 25{\cdot}8^\circ$ $47 \pm 18^\circ$	$(0{\cdot}02 \pm 0{\cdot}0067)\,W$ $(0{\cdot}0245 \pm 0{\cdot}0029)\,W$
Wu (1968)	Laboratory	0	$0{\cdot}048\,W$
Smith (1968)	Surface oil movement	$1 \pm 15^\circ$	$(0{\cdot}0346 \pm 0{\cdot}0071)\,W$
Dobrokonskii & Lesnikov (1972)	Laboratory	0	$(0{\cdot}0195 \pm 0{\cdot}0015)\,W$
Csanady (1972)	Lake Huron	$31{\cdot}7 \pm 57^\circ$	$(0{\cdot}07 \pm 0{\cdot}05)\,W$
Shemdin (1972)	Laboratory	0	$0{\cdot}03\,W$
Wright & Keller (1971)	Laboratory	0	$(0{\cdot}045 \pm 0{\cdot}014)\,W$
Kullenburg (1976)			$(0{\cdot}018\text{–}0{\cdot}036)\,W$
Wu (1975)	Laboratory	0	$0{\cdot}53\,u_* a$

TABLE 2. Field and laboratory data on drift currents.

The value of E will increase and the motion will be in the Ekman zone. The thickness of the layer dominated by the Ekman mechanism will be of the order of a wavelength of the main energy-containing waves, by virtue of the fact that $E = O(1)$. Beneath the Ekman layer the turbulent intensity will decay further, which will make E much larger than an order of one. The motion will essentially be an inviscid, inertial motion.

Under a given wind condition, any one or a combination of these three modes of motions will be possible, depending upon the value of ν_e and its vertically variation in magnitude. Thus a crucial question has to be asked: is the classical pure Ekman drift a true model for the surface drift current? This question is perhaps still best answered by the statement given by Ekman (1953) himself, after extensive field work to prove the existence of such flow, which is quoted as follows:

> The final result of the investigation may be summed up by saying that the observations made are not sufficient to establish definitely the existence of a 'pure drift current' as demanded by the theory, but that they are consistent – and in some respect even show a remarkable agreement – with the theoretical characteristics of such drift currents.

The summary of similar studies given in table 2 reflect the same results.

New parameters have to be incorporated to make a complete solution realistic. Nevertheless, the classical Ekman model is still regarded as the total solution in most ocean circulation studies.

Using the generalized Ekman equation (19) and figure 2, it is clear that the classical Ekman solution is one of the special cases within a whole range of solutions. Specifically the classical Ekman solution requires that the eddy viscosity be a constant and that the motion be strictly two-dimensional and dependent on z only. Leibovich (1977) demonstrated conclusively that vertical motions exist in the top layer of the ocean and such motions are part of the main mechanism generating and maintaining the Langmuir circulation. The existence of the Langmuir circulation is obvious to the most casual observer, but no consideration of this motion is included in the classical Ekman model.

The generalized Ekman equation presented here does contain the mechanisms necessary to model the whole range of solution. The different modes of motion can be combined to explain various phenomena in a unified way. Surface waves contribute in two ways: through a contribution from the low frequency waves to the Stokes drift and a contribution from the high frequency waves in determining the eddy viscosity ν_e. The interplay of waves and wind stresses can produce Langmuir cells superimposed on an Ekman spiral, as observed by Ichiye (1967). The importance of including both wind and waves when modelling the surface drift in the future is clear.

5. Conclusion

A model of surface-layer drift currents is proposed. It can be seen from the generalized Ekman equation and the classification chart of the surface drift current in figure 2 that the surface drift current may be a combination of three basic modes. This model of surface-layer dynamics is shown schematically in figure 3.

In this new model, the surface layer is most probably controlled by the frictional force, and hence is a Langmuir layer. The thickness of this layer is up to the depth of direct influence of the breaking waves, i.e. of the order of the amplitude of the main

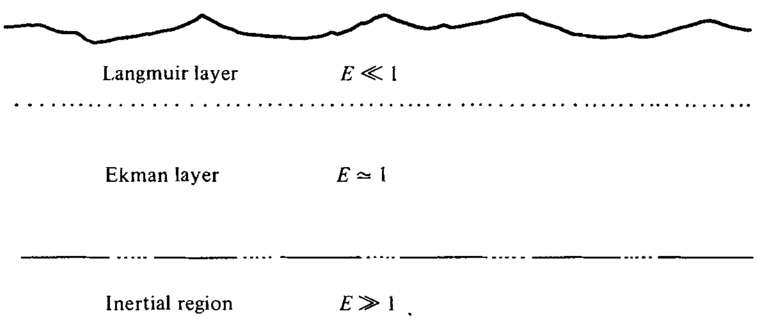

FIGURE 3. A proposed model of surface-layer drift currents.

energy-containing waves. Within this layer the turbulent intensity is high, and $E \ll 1$. The next layer, which may coexist with the Langmuir layer, is the Ekman layer, where the Coriolis force and the viscous force are of equal importance. This does not imply that the flow is of the same form as the classical Ekman solution. Vertical motions may still play a important role. It is called an Ekman layer only because of the balance between the Coriolis and the frictional forces. Depending upon the local state of the sea, and hence the magnitude and vertical variation of the eddy viscosity, the Langmuir layer may disappear or coexist with the Ekman layer in a complicated way. An inertial layer will exist below the region of influence of active wave motion and wind stresses. The fact that classical Ekman drift has not been observed consistently can be explained by this modified surface-layer model. Most wind-generated ocean circulation models (see, for example, Stern 1975) accept the Ekman drift model as the interface mechanism. This mechanism needs to be re-examined, particularly in light of Leibovich's major breakthrough. A unified method incorporating the contributions from both the wind and the wave motions is proposed in this model of surface-layer drift currents.

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