THEORETICAL INTERPRETATION OF FETCH LIMITED WIND–DRIVEN SEA OBSERVATIONS

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1. INTRODUCTION

Among field studies of wind-generated surface gravity waves on deep water, the case of a fetch limited wave growth is of special importance. This case is characterized by a constant (in space and time) offshore wind at a normal angle to the straight coastline. Under these idealized conditions, the wave field depends only on the fetch x which is defined as a distance from the shoreline. Highly interesting data sets of field observations are available, for example, from the following studies.

1. Measurements in Nakata Bay by Mitsuyasu *et al* (1971).
2. The JONSWAP studies by Hasselmann *et al* (1973).
3. Measurements in the Bothnian sea by Kahma (1981).
4. Measurements in Lake Ontario by Donelan *et al* (1985).
5. Measurements off the Nova Scotia coast by Dobson *et al* (1989).
6. Measurements in Lake St. Clair, Canada, by Donelan *et al* (1992).

The results of these experiments were analyzed by many authors and summarized by I. Young (1999).

In the present article we suggest a theoretical explanation of key results obtained in these observations. We show that most observations under limited fetch conditions find a satisfactory explanation within a rather simple theoretical framework based on Hasselmann's non-uniform stationary kinetic equation for spectral density of wave action in the absence of external forcing and dissipation. In particular, the spatial evolution of a spectral peak - at least for moderate values of the fetch - is in a good quantitative agreement with a similarity solution to this equation.

This statement may appear counter to physical intuition because surface waves are caused by wind and accompanied with a substantial energy dissipation due to wave breaking. This apparent controversy is resolved by noticing that the similarity solution mentioned above depends on two arbitrary constants. Their values are determined by the spectrum behavior at high frequencies. This spectrum tail is controlled by a complicated interplay of three major factors of wave dynamics in the short-fetch regime: the energy input from wind, energy dissipation due to white capping, and spectral energy transfer due to four–wave interaction. Determination of these unknown fundamental constants is a difficult problem, which is well outside the scope of this short article.

At larger values of the wind fetch, appropriate for open ocean conditions when the advection of the wave action due to wave group velocity becomes an important factor and the spectral transfer is controlled to a large extent by the inverse cascade of wave action, similarity solutions also exist. This case was treated by Zakharov and Zaslvaskii (1983) and Glazman (1994) yielding relationships that complement our present results.

Let us emphasize that the present theoretical model is based on the "first principles" and does not include any adjustable parameters. Nevertheless, the model yields results which are not only qualitative correct but are in a good quantitative agreement with field observations.

2. BASIC EQUATIONS

Let $\eta(\vec{r},t)$, $\vec{r}=(x,y)$ be the surface elevation field, $\eta(k,t)$ its Fourier transform, and $I_k = I_{-k} =< |\eta_k|^2 >$ the spectral density as a function of wave number. The wave field is described by the wave action spectral density N_k:

$$I_k = \frac{\omega_k}{2}(N_k + N_{-k}). \tag{1}$$

Here $\omega_k = \sqrt{gk}$ is the dispersion relation, where g is the vertical acceleration due to gravity.

The action density $N_k(\vec{r},t)$ satisfies the Hasselmann kinetic equation,

$$\frac{\partial N}{\partial t} + \frac{\partial \omega}{\partial \vec{k}}\frac{\partial N}{\partial \vec{r}} = S_{nl} + S_F, \tag{2}$$

where

$$\begin{aligned} S_{nl} &= \pi\, g^2 \int |T_{kk_1k_2k_3}|^2\, \delta(k+k_1-k_2-k_3)\, \delta(\omega_k+\omega_{k_1}-\omega_{k_2}-\omega_{k_3}) \times \\ & (N_{k_1}N_{k_2}N_{k_3} + N_kN_{k_2}N_{k_3} - N_kN_{k_1}N_{k_2} - N_kN_{k_1}N_{k_3})\, dk_1\, dk_2\, dk_3 \end{aligned} \tag{3}$$

is a nonlinear interaction term and $T_{kk_1k_2k_3}$ is a coupling coefficient. The exact expression for T can be found in (Webb 1973, Zakharov 1999). Of great importance is the fact that $T_{kk_1k_2k_3}$ is a homogenous function of the third order:

$$T_{\zeta k,\zeta k_1,\zeta k_2,\zeta k_3} = \zeta^3\, T_{kk_1k_2k_3}. \tag{4}$$

The term $S_F = S_{in} + S_{ds}$ represents a source function which includes the wind input S_{in} and the breaking wave induced dissipation S_{ds}.

By definition, the variance of the surface elevation is given by

$$\sigma^2 = \int \omega_k\, N_k\, dk, \tag{5}$$

and the mean frequency is

$$\bar{\omega} = \frac{1}{\sigma^2}\int \omega_k^2\, N_k\, dk. \tag{6}$$

The overall wave steepness is characterized by an integral quantity

$$\hat{\alpha} = \frac{\sigma^2\,\bar{\omega}^4}{g^2}. \tag{7}$$

Let U_a be the wind speed at a reference height a. We can now introduce the characteristic angular frequency and wave number related to the given wind speed U_a by

$$\omega_0 = g/U_a, \quad k_0 = g/U_a^2$$

Most authors agree that S_{in} can be presented in the form

$$S_{in} = \mu\, F(\xi)\, \omega\, N(k), \tag{8}$$

where $\xi = \omega\, \cos\theta/\omega_0$ and

$$\mu = 0.1 \sim 0.3\, \frac{\rho_a}{\rho_\omega} \simeq 10^{-4}. \tag{9}$$

There is no agreement about the exact form of the function $F(\xi)$. According to Snyder *et al* (1983),

$$F(\xi) = \begin{cases} \xi - 1 & \xi > 1 \\ 0 & \xi < 1. \end{cases} \tag{10}$$

According to Hsiao and Shemdin (1983), also Donelan *et al* (1984),

$$F(\xi) = \begin{cases} (\xi - 1)^2 & \xi > 1 \\ 0 & \xi < 1. \end{cases} \tag{11}$$

Tolman and Chalikov (1996) proposed a more complicated form of $F(\xi)$. In any case, all such models have $F(\xi) \simeq 1$ for $\xi \sim 1$.

An analytic expression for S_{ds} is much less certain. Komen *et al* (1984) proposed the following form:

$$S_{ds} = -3.33 \times 10^{-5} \left(\frac{\hat{\alpha}}{\alpha_{pm}}\right)^4 \left(\frac{\omega}{\bar{\omega}}\right)^2 \omega\, N, \tag{12}$$

where $\alpha_{pm} = 4.57 \times 10^{-3}$.

This relationship is presently used in WAM and SWAN wave prediction models. In our opinion, expression (12) overestimates S_{ds} in the spectral range $\omega_0 < \omega < 3 \sim 4\,\omega_0$. The wave breaking is insignificant in the spectral peak range. It becomes important only at high frequencies and wavenumbers. Therefore, it is reasonable to assume that S_{ds} is small near the peak and rapidly grows with an increasing frequency at $\omega \gg \omega_0$. In the vicinity of the spectral peak - which is of our main interest - one can write:

$$S_f \simeq S_{in} \simeq \mu\, F(\xi)\, \omega\, N(k). \tag{13}$$

3. SELF–SIMILAR SOLUTION

Let us introduce dimensionless variables :

$$x = \chi/k_0, \quad \vec{k} = k_0\, \vec{\kappa}, \quad \omega = \omega_0\, \Omega, \quad \Omega = \sqrt{\kappa},$$

yielding

$$N(k) = \frac{1}{\omega_0\, k_0^4}\, n(\vec{\kappa}). \tag{14}$$

The non–dimensional surface height variance and the non–dimensional average frequency are:

$$\epsilon = k_0^2\, \sigma^2 = \frac{g^2}{U_0^4}\, \sigma^2, \tag{15}$$

$$\nu = \frac{1}{2\pi}\, \frac{\bar{\omega}}{\omega_0}. \tag{16}$$

Both ϵ and ν can be expressed in terms of $n(\vec{\kappa})$. Apparently,

$$\epsilon = \int \sqrt{\kappa}\, n(\vec{\kappa})\, d\vec{\kappa}, \tag{17}$$

$$\nu = \frac{1}{2\pi}\, \frac{\int |\kappa|\, n(\kappa)\, d\vec{\kappa}}{\int \sqrt{|\kappa|}\, n(\kappa)\, d\kappa}. \tag{18}$$

Under limited-fetch conditions, $n(\vec{\kappa}, \chi)$ is governed by the kinetic equation

$$\frac{\cos\theta}{2\Omega}\, \frac{\partial n}{\partial \chi} = \tilde{S}_{nl} + \tilde{S}_F, \tag{19}$$

where

$$\begin{aligned} \tilde{S}_{nl} &= \pi \int |T_{\kappa\kappa_1\kappa_2\kappa_3}|^2\, \delta(\vec{\kappa} + \vec{\kappa_1} - \vec{\kappa_2} - \vec{\kappa_3})\, \delta(\Omega + \Omega_1 - \Omega_2 - \Omega_3) \times \\ &\quad (n_{\kappa_1} n_{\kappa_2} n_{\kappa_3} + n_\kappa n_{\kappa_2} n_{\kappa_3} - n_\kappa n_{\kappa_1} n_{\kappa_2} - n_\kappa n_{\kappa_1} n_{\kappa_3})\, d\kappa_1\, d\kappa_2\, d\kappa_3 \sim |\kappa|^{19/2}\, n^3, \\ \tilde{S}_F &\simeq \mu\, \Omega\, F(\Omega)\, n(\kappa, \chi). \end{aligned} \tag{20}$$

In (19), $\mu \simeq 10^{-4}$ is a small parameter. In a first approximation, we can set $\mu = 0$, and obtain the "conservative" kinetic equation:

$$\frac{\cos\theta}{2\Omega}\,\frac{\partial n}{\partial \chi} = \tilde{S}_{nl}. \tag{21}$$

This governing equation is the main focus of our analytical effort. It contains a family of self–similar solutions. In polar coordinates κ, θ on $\vec{\kappa}$ plane these solutions can be presented as:

$$n(\kappa,\theta,\chi) = a\,\chi^{\alpha}\,P_\beta(b\,\chi^{\beta}\,\kappa,\theta). \tag{22}$$

Here, a, b, α, β are constants.

Substituting (22) into (21) one finds

$$\alpha = 5\,\beta - 1/2, \quad a = b^5. \tag{23}$$

Ultimately,

$$n(\kappa,\theta,\chi) = b^5\,\chi^{5\beta-1/2}\,P_\beta(b\,\chi^{\beta}\,\kappa,\,\theta). \tag{24}$$

In (24), β and b are yet unknown constants, and $P_\beta(z,\theta)$ is a function of two variables with $z = b\chi^\beta\kappa$. Let us emphasize that this function is independent of b and it satisfies the following integro–differential equation:

$$\begin{aligned}
&\frac{\cos\theta}{2\sqrt{z}}\,[(5\beta-1/2)P_\beta + \beta z P_z] = \pi\int |T_{z,z_1,z_2,z_3,\theta,\theta_1,\theta_2,\theta_3}|^2 \times \\
&\delta(z\cos\theta + z_1\cos\theta_1 - z_2\cos\theta_2 - z_3\cos\theta_3)\,\delta(z\sin\theta + z_1\sin\theta_1 - z_2\sin\theta_2 - z_3\sin\theta_3) \times \\
&\delta(\sqrt{z} + \sqrt{z}_1 - \sqrt{z}_2 - \sqrt{z}_3)\Big[P_\beta(z_1,\theta_1)\,P_\beta(z_2,\theta_2)\,P_\beta(z_3,\theta_3) + \\
&P_\beta(z,\theta)\,P_\beta(z_2,\theta_2)\,P_\beta(z_3,\theta_3) - P_\beta(z,\theta)\,P_\beta(z_1,\theta_1)\,P_\beta(z_2,\theta_2) - \\
&P_\beta(z,\theta)\,P_\beta(z_1,\theta_1)\,P_\beta(z_3,\theta_3)\Big] z_1\,z_2\,z_3\,dz_1\,dz_2, dz_3\,d\theta_1\,d\theta_2\,d\theta_3.
\end{aligned} \tag{25}$$

This equation has to be solved numerically. Let us denote:

$$A_\beta = \int \sqrt{z}\,P_\beta(z,\theta)\,z\,dz\,d\theta, \tag{26}$$

$$B_\beta = \int z\,P_\beta(z,\theta)\,z\,dz\,d\theta. \tag{27}$$

From (17), (18) one finds

$$\begin{aligned}
\epsilon &= b^{5/2}\,\chi^{\frac{5\beta-1}{2}}\,A_\beta, \\
\nu &= \frac{1}{2\pi}\,b^{-1/2}\,\chi^{-\beta/2}\,\frac{B_\beta}{A_\beta}.
\end{aligned} \tag{28}$$

These equations can be written as:

$$\begin{aligned}
\epsilon &= u\,\chi^{p}, \\
\nu &= v\,\chi^{-q}\ ,
\end{aligned} \tag{29}$$

where

$$q = \frac{2p+1}{10}, \qquad \beta = \frac{2p+1}{5}, \tag{30}$$

$$v = \frac{1}{2\pi}\,\bar{u}^{1/5}\,C_\beta, \qquad C_\beta = \frac{B_\beta}{A_\beta^{1/5}}. \tag{31}$$

Integrating equation (19) over $\vec{\kappa}$ one obtains the balance equation

$$\frac{1}{2}\int \frac{\cos\theta}{\Omega}\,\frac{\partial n}{\partial \chi}\,d\vec{\kappa} = \int \tilde{S}_F\,d\kappa = Q(\chi). \tag{32}$$

Here Q is the total input of wave action - a net result of wind forcing and breaking wave dissipation. Substituting the self–similar solution (24) into (25), one finds

$$Q \simeq \chi^{\frac{7\beta-3}{2}}. \tag{33}$$

Therefore, the solution to (24) implies the presence of a wave action source Q at high wave numbers. For $\beta = \beta_{crit} = 3/7$, the net input $Q = const$, and the intensity of the source does not depend on the fetch. For $\beta > \beta_{crit}$, the input Q grows with an increasing fetch. In the presence of $\tilde{S}_F$, the self–similar solution is valid only for χ which are not too large. For $\Omega \sim 1$, the left hand side of (19) can be estimated as n/χ, while $\tilde{S}_F \simeq \mu n$. Thus, the self–similar solution remains valid for

$$\chi < \frac{1}{\mu} < 10^4. \tag{34}$$

Note, that a special self–similar solution

$$\begin{aligned} &n_{crit}(k,\theta,\chi) = b\,\chi^{23/14}\,P_{3/7}\,(b\,\chi^{3/7}\,\kappa,\theta),\\ &p = \frac{n}{7} = 0.57,\\ &q = \frac{3}{14} = 0.21, \end{aligned} \tag{35}$$

corresponding to constant net input $Q = const$, was studied in papers of Zakharov and Zaslavskii (1982, 1983) and Glazman (1994).

4. COMPARISON WITH EXPERIMENT

We shall now compare the similarity solution (24) with the field observations under limited-fetch conditions. Let us first notice that an elementary analysis of observed spectra shows their self–similar behavior. The similarity is implicit in the fact that the spectra can be expressed by a universal form that involves a finite number of parameters.

The frequency spectrum can be introduced as follows:

$$F(f) = k\,\omega_k\,\frac{dk}{df}\int_0^{2\pi} N(k,\theta)\,d\theta, \quad k = (2\pi\,f)^2/g. \tag{36}$$

For example, the results of the JONSWAP experiment are summarized by the spectral form (Hasselmann *et al*, 1973):

$$F(f) = \alpha\,g^2\,(2\pi)^{-4}\,f^{-5}\,\exp\left[-3/4\left(\frac{f}{f_p}\right)^{-4}\right]\gamma^{\exp\left[-\frac{(f-f_p)^2}{2\sigma^2\,f_p^2}\right]} \tag{37}$$

At moderate values of the non–dimensional fetch $\chi < 10^5$, parameters γ and σ are approximately constant

$$\gamma \simeq 3.3, \quad \sigma \simeq 0.08,$$

while α and f_p are powers of χ. Form (33) is explicitly self–similar. Actually, the f^{-5} asymptotic of (37) represents the regime described by the Phillip spectrum (1958). A more accurate analysis of the JONSWAP spectra (J. Battjes, 1985) shows that a f^{-4} behavior is more relevant. Donelan *et al* (1984) also found the f^{-4} asymptotic for their spectra. This is exactly what the theory predicts, but a discussion of this issue is outside the scope of this article.

In accordance with the main assumption of self-similarity, the fetch dependence of ϵ and ν is described by powerlike functions of χ. To assure that the spectra are described by the similar solution (24), one has to compare these functions with the theoretically predicted forms (29). The results of the major fetch limited experiments can be summarized in two tables:

Table 1

Study	$\epsilon(\chi)$	$\nu(\chi)$
Nakata Bay, Mitsuyasu *et al* (1971)	$\epsilon = 2.89 \times 10^{-7}\,\chi$	$\nu = 3.12\,\chi^{-0.33}$
JONSWAP, Hasselman *et al* (1973)	$\epsilon = 1.6 \times 10^{-7}\,\chi$	$\nu = 3.5\,\chi^{-0.33}$
Bothnian Sea, Kahma (1981)	$\epsilon = 2.6 \times 10^{-7}\,\chi$	$\nu = 3.18\,\chi^{-0.33}$
Lake St. Clair, Donelan *et al* (1992)	$\epsilon = 1.7 \times 10^{-7}\,\chi$	$\nu = 3.6\,\chi^{-0.33}$

Table 2

Study	$\epsilon(\chi)$	$\nu(\chi)$
Lake Ontario, Donelan *et al* (1985)	$\epsilon = 8.415 \times 10^{-7}\,\chi^{0.76}$	$\nu = 1.85\,\chi^{-0.23}$
North Atlantic, Dobson *et al* (1989)	$\epsilon = 12.7 \times 10^{-7}\,\chi^{0.75}$	$\nu = 1.7\,\chi^{-0.24}$

Expression (30) makes it possible to find q if p is give. For the first group of experiments $p = 1$, and one obtains:

$$q = 0.3, \quad \beta = 0.6$$

For the Lake Ontario experiment:

$$p = 0.76, \quad q = 0.25, \quad \beta = 0.5$$

For the North Atlantic study:

$$p = 0.75, \quad q = 0.25, \quad \beta = 0.5$$

In (31), coefficients C_β are unknown constants defined by the shapes of the solutions to equation (25). Let us denote:

$$C|_{\beta=0.6} = C_1, \quad C|_{\beta=0.5} = C_2$$

Now we can compare experimental and theoretical results for $\nu(\chi)$. These results are summarized in Table 3.

Table 3

Study	**Experiment**	**Theory**	**Optimized theory**
Nakata Bay	$3.12\,\chi^{-0.33}$	$3.23\,C_1\,\chi^{-0.30}$	$3.20\,\chi^{-0.30}$
JONSWAP	$3.5\,\chi^{-0.33}$	$3.64\,C_1\,\chi^{-0.30}$	$3.6\,\chi^{-0.30}$
Bothnian Sea	$3.18\,\chi^{-0.33}$	$3.1\,C_1\,\chi^{-0.30}$	$3.06\,\chi^{-0.30}$
Lake St. Clair	$3.6\,\chi^{-0.33}$	$3.6\,C_1\,\chi^{-0.30}$	$3.56\,\chi^{-0.30}$
Lake Ontario	$1.85\,\chi^{-0.23}$	$2.6\,C_2\,\chi^{-0.25}$	$1.84\,\chi^{-0.25}$
North Atlantic	$1.7\,\chi^{-0.24}$	$2.4\,C_2\,\chi^{-0.25}$	$1.7\,\chi^{-0.25}$

Coefficients C_1 and C_2 in the second column of Table 3 are not adjustable parameters. They have definite values to be found by numerical solution of equation (25). We plan to determine these values in a later study. At the present time we propose a hypothesis that their values are "optimal", so that the third column in Table 3 is sufficiently close to the first column. Optimization by the least square method gives:

$$C_1 = 0.99, \quad C_2 = 0.71$$

Table 3 demonstrates a good agreement between theory and experiment.

5. DISCUSSION

1. The close agreement between the theoretical and experimental results indicates that the Hasselmann kinetic equation is an adequate model describing the evolution of wind–driven surface gravity waves. Moreover, the evolution of the spectral peak at moderate fetches can be faithfully described by the "conservative" kinetic equation where the forcing and dissipation terms are dropped.

Self–similar solution (24) describes downshift of the peak frequency. This is a direct consequence of the "inverse cascade" of energy and wave action. The physical origin of the inverse cascade is the existence of an additional integral of motion - the wave action. The wave action is preserved only in four–wave interactions. One can say that the very fact of the downshifting of the spectral peak indicates a dominant role of four–wave nonlinear interaction.

This qualitative analysis finds, as we just showed, a reasonable quantitative confirmation.

2. Different groups of experiments give two different values for β: $\beta = 0.6$ and $\beta = 0.5$. Both of them are larger than the critical value $\beta = 3/7 \simeq 0.43$. This means that the input of wave action increases with an increasing fetch. The explanation here is rather simple. According to (7), (29), the characteristic steepness $\hat{\alpha}$ decreases as the fetch increases:

$$\hat{\alpha} \simeq \chi^{\frac{\beta-1}{2}}. \tag{38}$$

According to the Komen's form (12) (and to many other models of the breaking wave dissipation), S_{ds} is very sensitive to $\hat{\alpha}$ and it rapidly decreases as $\hat{\alpha}$ decreases. As a result, the white capping is more vigorous for "young" seas. When χ grows, S_{ds} becomes suppressed and the wave action input from wind increases.

3. According to the theoretical prediction, the self–similar solution (24) is valid only for not too large fetches. What happens after? It depends on the structure of S_F for long waves. If, as it was estimated in models (11, 12), $S_F = 0$ for $\omega < \omega_0$, the downshift continues infinitely long. Asymptotically it is described by the self–similar solution with $\beta = \beta_{crit} = 3/7$. Behind the spectral peak, the spectrum has an asymptotic form $F(f) \simeq f^{-11/3}$.

In the experiments of Donelan *et al* (1992) performed on Lake St. Clair, as well as in the earlier experiments of Pierson and Moskowitz (1964), and SMB CERC (1977), the downshift is arrested approximately at $\chi_{crit} \simeq 5 \times 10^4$. It corresponds to a very long fetch, $\chi_{crit} \simeq 10^4 L$, where L is a characteristic wave length. In a typical case $L \simeq 100\, m$, thus $\chi_{crit} \simeq 10^3 km$.

According to J. Young, who summarized these results, stabilization of the downshift (inverse cascade) is going up to the level:

$$\epsilon = 4 \sim 5 \times 10^{-3}, \quad \nu \simeq 0.13$$

Thus, $\bar{\omega} \simeq 0.81\, \omega_0$. In other words, the maximal "wave age", observed in these experiments, does not exceed unit. In this scenario, when $\chi \to \chi_{crit}$, the self–similarity is violated. γ is not a constant any more. At $\chi \to \chi_{crit}$, $\gamma \to 1$, and the spectrum loses its conspicuous peak. This process, "sea maturing", was observed by Pierson and Moskovitz in 1964, and by Donelan *et al* in 1992.

Not all researchers agree with this concept. According to Glazman (1994), the inverse cascade is not arrested at wave age of order of one, but continue to the spectral area of waves with mean frequency $\omega < \omega_0$. In his experiments near Hawaii island, he observed the wave of age $\omega_0/\omega \geq 3$. Anyway, downshift in this area is more slow than predicts the critical self–similar solution (35), and the spectrum tail becomes less steep: $F(f) \simeq f^{-3}$. This question urgently needs more experimental studies.

Nevertheless, it is clear that both the slowdown and the arrest of the inverse cascade ("maturing of the sea"), occur due to dissipation of long waves with phase velocities close to the wind speed or exceeding it. The decrement of this dissipation, β, is very small: $\beta/\omega \simeq \mu \sim 10^{-4}$. A mechanism of this dissipation is still unclear. It could be a combination of wave breaking and friction over a turbulent air. Meanwhile, a scenario of the inverse cascade, the slowdown and arrest is very sensitive to the exact value and the details of this dissipation. It makes the problem of wave prediction extremely difficult from a theoretical viewpoint.

4. The ideas presented in this article can be applied to wind–driven waves on a finite–depth fluid. However, this problem is much more difficult. On a finite–depth the kinetic equation has now self–similar solutions and the solutions for comparison with experiment can be obtained only by the use of a massive numerical simulation. Anyway, the idea of predominance of four–wave nonlinear interaction in the area of spectral peak is still applicable, at least as a first approximation. The next important factor to be taken into consideration is a bottom friction.

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