Dephasing and delay time fluctuations in the chaotic scattering of a quantum particle weakly coupled to a complicated background

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Abstract

Effect of a complicated many-body environment is analyzed on the chaotic motion of a quantum particle in a mesoscopic ballistic structure. The dephasing and absorption phenomena are treated on the same footing in the framework of a schematic microscopic model. The single-particle doorway resonance states excited in the structure via an external channel are damped not only because of the escape onto such channels but also due to ulterior population of the long-lived background states. The transmission through the structure is presented as an incoherent sum of the flow formed by the interfering damped doorway resonances and the retarded flow of the particles reemitted by the environment. The resulting internal damping as well as the dephasing rate are uniquely expressed in terms of the spreading width which controls the coupling to the background. The formation of the long-lived fine-structure resonances strongly enhances delay time fluctuations thus broadening the delay time distribution.

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1. Analog experiments with open electromagnetic microwave cavities \[1, 2, 3, 4\] on the one hand and extensive study of the electron transport through ballistic microstructures \[5\] on the other have drawn during last decade much attention to peculiarities of chaotic wave interference in open billiard-like set-ups. It is well recognized by now that the statistical approach \[6, 7, 8\] based on the random matrix theory (RMT) provides a reliable basis for describing the universal fluctuations characteristic of such an interference which, in particular, manifests itself in the single-particle resonance chaotic scattering and transport phenomena. At the same time, experiments with ballistic quantum dots reveal persisting up to very low temperatures appreciable deviations from the predictions of RMT which indicate some loss of the quantum coherence.

Two different methods of accounting for the dephasing effect have been suggested which give different results. In the Büttiker’s voltage-probe model \[9\] an subsidiary randomizing scatterer is introduced with \(M_\phi\) channels each with a sticking coefficient \(P_\phi\). Even assuming all these channels to be equivalent we are still left with two independent parameters. This results in an ambiguity since quantities of physical interest (e.g. the conductance distribution) depend, generally, on \(M_\phi\) and \(P_\phi\) separately whereas the dephasing phenomenon is controlled by the unique parameter: the dephasing time \(\tau_\phi\) which is fixed by their product. On the other hand, only the strength of a uniform imaginary potential governs the Efetov’s model \[10\] but, implying absorption, this model suffers a loss of unitarity. The deficiencies and differences of the two models \[11, 12\] have been analyzed, in particular in \[12\]. A prescription was suggested how to get rid of uncertainties, and simultaneously, to accord the models by considering the limit \(M_\phi \to \infty, P_\phi \to 0\) at fixed product \(M_\phi P_\phi \equiv \Gamma_\phi\) in the first model and by compulsory restoration of the unitarity in the second. The construction proposed infers a complicated internal structure of the probe. Otherwise, the assumed limit could hardly be physically justified. If so, the typical time \(\tau_p\) spent by the scattering particle inside the probe, being proportional to its mean spectral density, forms, generally, a new time scale different from the dephasing time. A question thus arises on the total time delay during transport through a ballistic mesoscopic structure in presence of a complicated background.

2. Below we propose a model of dephasing and absorption phenomena, which from the very beginning does not suffer any ambiguity. We consider the environment as a complicated many-body system with a very dense energy spectrum. The evolution of the extended
system: the moving inside the structure particle and the background coupled to each other with an interaction $V$, is described by means of an enlarged non-Hermitian Hamiltonian matrix $\mathcal{H}$ of order $N = N^{(s)} + N^{(e)}$

$$\mathcal{H} = \begin{pmatrix} \mathcal{H}^{(s)} & V^\dagger \\ V & \mathcal{H}^{(e)} \end{pmatrix} \quad (1)$$

The two blocks $\mathcal{H}^{(r)} = H^{(r)} + \frac{i}{2} A^{(r)} A^{(r)\dagger}$, $r = s, e$ ($N^{(e)} \gg N^{(s)}$) along the main diagonal represent the non-Hermitian effective Hamiltonians of the bare single-particle open system and the environment respectively. The Hermitian matrices $H^{(r)}$ describe their closed counterparts (the corresponding mean level spacings satisfy the condition $D(r = s) \gg d(r = e)$) when the matrices $A^{(r)}$ are built of the amplitudes connecting the internal and channel states. The two (generally unstable) subsystems are mixed by the off-diagonal $N^{(e)} \times N^{(s)}$ coupling matrix $V$. When uncoupled, these subsystems have no common decay channels so that the coupling $V$ is purely Hermitian. This means that the background states have no direct access to the observable outer channels and attain it only due to the mixing with the single-particle ”doorway” states inside the dot.

The total $M \times M$ scattering matrix $S(E) = I - iA^\dagger (E - \mathcal{H})^{-1} A$ is unitary so far as all the $M = M^{(s)} + M^{(e)}$ open channels, observable ($s$) as well as hidden ($e$), are taken into account. However, only transitions between $M^{(s)}$ outer channels are accessible for direct observations. They are described by the $M^{(s)} \times M^{(s)}$ submatrix $S(E) = I - i\mathcal{T}(E)$ where

$$\mathcal{T}(E) = A^{(s)\dagger} \mathcal{G}^{(s)}_D(E) A^{(s)} \quad (2)$$

and the matrix

$$\mathcal{G}^{(s)}_D(E) = \frac{1}{E - \mathcal{H}^{(s)} - \Sigma^{(s)}(E)} \quad (3)$$

is the upper left $N^{(s)} \times N^{(s)}$ block of the resolvent $\mathcal{G}(E) = (E - \mathcal{H})^{-1}$. The subscript $D$ stands in (3) for ”doorway”. The self-energy matrix

$$\Sigma^{(s)}(E) = V^\dagger \frac{1}{E - \mathcal{H}^{(e)}} V \equiv V^\dagger \mathcal{G}^{(e)}_0(E) V \quad (4)$$

includes all virtual transitions between the bare doorway and environment resonances. This matrix is not Hermitian as long as the background states are not stable. Therefore the submatrix $S$ is not, generally, unitary. The flow of outgoing particles through outer channels reduces because of inelastic processes in the background thus implicating absorption.
We further assume that the coupling matrix elements $V_{\mu m}$, where the indices $\mu$ and $m$ span the background’s and the system’s Hilbert spaces respectively, are random Gaussian variables, $\langle V_{\mu m} \rangle = 0$; $\langle V_{\mu m} V_{\nu n}^* \rangle = \delta_{\mu \nu} \delta_{mn} \frac{1}{2} \Gamma_s d$. Here $\Gamma_s = 2\pi \langle |V|^2 \rangle / d$ is the spreading width \[13, 14, 15\]. The inequality $\langle |V|^2 \rangle \gg d^2$ is implied so that the interaction $V$, though weak, is not weak enough for perturbation theory to be valid. Retaining the original notations also for the quantities averaged over the interaction $V$, we find in the main $1/N^{(e)}$ approximation

\[ G^{(s)}_b(E) = \left[ E - \frac{1}{2} \Gamma_s g^{(e)}(E) - \mathcal{H}^{(s)} \right]^{-1}. \]

where $g^{(e)}(E) = \frac{2}{\pi} \text{Tr} G_0^{(e)}(E)$. The resonance spectrum $\{ \mathcal{E}_a = E_a - \frac{i}{2} \Gamma_a \}$ is defined in such an approximation from $N^{(s)}$ similar equations

\[ \mathcal{E} - \frac{1}{2} \Gamma_s g^{(e)}(\mathcal{E}) - \mathcal{H}^{(s)} = 0 \]

originated each from one of the bare single-particle doorway resonances.

In what follows we analyze in detail temporal aspects of the scattering where the influence of the background shows up in the most full and transparent way. In particular, a straightforward calculation gives for the Smith delay-time submatrix $Q = -iS^dS/dE$ in the $M^{(s)}$-dimensional subspace of the observable channels $Q = b^{(s)}b^{(s)} + b^{(e)}b^{(e)} = Q^{(s)} + Q^{(e)}$ where the matrix amplitudes

\[ b^{(s)}(E) = G^{(s)}(E)A^{(s)}; \quad b^{(e)}(E) = G^{(e)}_0(E)A^{(e)}; \quad b^{(e)}(E) = G^{(e)}_0(E)V b^{(s)}(E) \]

have dimensions $N \times M^{(s)}, N^{(s)} \times M^{(s)}$ and $N^{(e)} \times M^{(s)}$ respectively. The two contributions $Q^{(s,e)}$ correspond to the modified due to the interaction with the background time delay within the dot and delay because of the virtual transitions into the background. After averaging over the interaction $V$ we arrive in the main approximation to

\[ Q(E) = \Lambda(E) b^{(s)}b^{(s)} = \Lambda(E)Q^{(s)}(E); \quad \Lambda(E) = 1 + \frac{1}{2} \Gamma_s l^{(e)}(E) \]

where $l^{(e)}(E) = \frac{2}{\pi} \text{Tr} \left[ G_0^{(e)}(E)G_0^{(e)}(E) \right]$. A given bare doorway resonance state with the complex energy $\mathcal{E}_0 = \varepsilon_0 - \frac{i}{2} \gamma_0$ generates a set of exact resonance states with complex energies which are found from the eq. (6). Since the energy spectrum of the environment is very dense and rich, its states are supposed to decay through a large number of weak statistically equivalent channels. The corresponding decay amplitudes $A^{(e)}_{\mu} \Delta$ are random Gaussian variables with zero means and variances
\[ \langle A_{\mu}^{(e)} A_{\nu}^{(e')} \rangle = \delta^{e e'} \delta_{\mu \nu} \gamma_e / M^{(e)}. \] The widths \( \gamma_e \) do not fluctuate when \( M^{(e)} \gg 1 \). The interaction \( V \) redistributes the initial widths over exact resonances as \( \Gamma_\alpha = f_\alpha \gamma_0 + (1 - f_\alpha) \gamma_e \) with the strength function \( f_\alpha \equiv f(E_\alpha) \) obeying the condition \( \sum_\alpha f_\alpha = 1 \).

Neglecting unobservable spectral fluctuations of the background we assume a rigid spectrum with equidistant levels \( \epsilon_\mu = \epsilon_0 + \mu d - \frac{i}{2} \gamma_e \). One of the advantages of this uniform model \( [13, 15] \) is that the loop functions \( g^{(e)}(E), l^{(e)}(E) \) as well as the strength function \( f_\alpha \) can be calculated explicitly \( [15] \). Depending on the magnitude of the coupling there exists two different scenarios. In the limit \( \Gamma_s \gg \gamma_0 - \gamma_e \) (the natural assumption \( \gamma_e \ll \gamma_0 \) is accepted throughout the paper) of strong interaction all the individual strengths \( f_\alpha \) are small, \( f(E_\alpha) \leq 2d / \pi \Gamma_s \), and are distributed around the energy \( \epsilon_0 \) according to the Lorentzian with the width \( \Gamma_s, f_\alpha \propto L_{\Gamma_s}(E_\alpha - \epsilon_0) \). Thus the original doorway state fully dissolves in the sea of the background states. In the opposite limit of weak coupling, \( \Gamma_s \ll \gamma_0 - \gamma_e \), which is the one of our interest, the strength \( f_0 = 1 - \Gamma_s / (\gamma_0 - \gamma_e) \) remains large when the rest of them are small again \( f(E_\alpha) \leq 2d \Gamma_s / \pi (\gamma_0 - \gamma_e)^2 \) and distributed as \( L_{(\gamma_0 - \gamma_e)}(E_\alpha - \epsilon_0) \). Therefore, only in the weak-coupling case the doorway state preserves individuality and can be observed through the outer channels. The interaction with the background surrounds a doorway resonance by a bunch of fine-structure resonances with Lorenzianly distributed heights.

3. When the scattering energy \( E \) approaches the environment ground energy only low-lying bare background states are involved which are stable, \( A^{(e)}_\mu = 0 \). The self-energy matrix \( \Sigma^{(s)}(E) \) is hermitian and, as a consequence, the scattering matrix \( S(E) \) is unitary in this limit. Note that the averaging over the coupling \( V \) does not spoil the unitarity.

\[ g^{(e)}(E) = \cot \left( \frac{E - \epsilon_0}{d} \right); \quad l^{(e)}(E) = \frac{\pi}{d} \sin^{-2} \left( \frac{E - \epsilon_0}{d} \right). \] (9)

If the structure is almost closed so that the considered resonance is isolated the individual cross sections

\[ \sigma_{ab}(E) = \frac{\gamma^{(a)} \gamma^{(b)}}{[E - 1/2 \Gamma_s \cot (\pi E/d)]^2 + 1/4 \gamma_0^2}; \quad \gamma^{(a)} = |A_0^{(a)}|^2; \quad \gamma_0 = \Sigma_\alpha \gamma^{(a)} \] (10)
as well as the Wigner time delay

\[ \tau_W(E) = \text{tr} Q(E) = \gamma_0 \frac{1 + (\pi \Gamma_s/2d) \sin^{-2} (\pi E/d)}{[E - 1/2 \Gamma_s \cot (\pi E/d)]^2 + 1/4 \gamma_0^2} \] (11)

(we have set \( \epsilon_0 = 0 \)) reveal strong fine-structure fluctuations on the scale of the background level spacing \( d \). In particular, the delay time fluctuates between \( \tau_W(E = \epsilon_\mu) = \)
\[ (2\pi/d)(\gamma_0/\Gamma_s) \sim 1/\Gamma_\mu \] at the points of the fine structure levels and a much smaller value \( \tau_W(E \approx 0 \neq \epsilon_\mu) \approx (2\pi/d)(\Gamma_s/\gamma_0) \) in between. In fact, a particular fine-structure resonance cannot be resolved and only quantities averaged over an energy interval \( d \ll \delta E \ll \Gamma_s \), \( D \) are observed. (Equivalently, one can use instead ensemble averaging over the fine-structure resonances.) Actually, it is enough because of periodicity to average over the interval \(-\frac{d}{2} < E < \frac{d}{2}\). The averaging which can easily be performed explicitly yields

\[
\sigma^{ab}(E) = \left(1 + \frac{\Gamma_s}{\gamma_0}\right) \frac{\gamma_a \gamma_b}{E^2 + \frac{1}{4}(\gamma_0 + \Gamma_s)^2} 
\] (12)

and

\[
\tau_W(E) = \frac{\gamma_0 + \Gamma_s}{E^2 + \frac{1}{4}(\gamma_0 + \Gamma_s)^2} + \frac{2\pi}{d}. 
\] (13)

The first terms in the both expressions correspond to excitation of a damped resonance with the total width \( \Gamma_0 = \gamma_0 + \Gamma_s \) when the second account for the particles reinjected from the environment after the typical time delay proportional to the background level density.

Generally, however, resonances overlap and the scattering amplitudes can be presented as a sum of a large number \( N^{(s)} \gg 1 \) of interfering resonance contributions each one being proportional to the energy dependent product

\[
\frac{1}{\mathcal{D}_n(E)\mathcal{D}^*_n(E)} = \frac{1}{\mathcal{E}_n - \mathcal{E}^*_n} \left[ \frac{1}{\mathcal{D}_n(E)} - \frac{1}{\mathcal{D}^*_n(E)} \right]
\] (14)

where \( \mathcal{D}_n(E) = E - \mathcal{E}_n - \frac{i}{2}\Gamma_s \cot\left(\frac{\pi(E - \mathcal{E}_n)}{d}\right) \) and \( n, n' = 1, 2, ..., N^{(s)} \). Using the identities:

\[
\frac{1}{\mathcal{D}_n(E)} = \frac{1}{E - \mathcal{E}_n + \frac{i}{2}\Gamma_s}; \quad \frac{1}{\mathcal{E}_n - \mathcal{E}^*_n} = i \int_0^\infty dt e^{-i(\mathcal{E}_n - \mathcal{E}^*_n) t}
\] (15)

we arrive finally at an incoherent sum \( \sigma^{ab}(E) = \sigma^{ab}_1(E) + \sigma^{ab}_2(E) \) of the flow formed by the interfering damped doorway resonances

\[
\sigma^{ab}_1(E) = \left| \left( \frac{A^\dagger}{E - \mathcal{H}^{(s)} + \frac{i}{2}\Gamma_s} \right)^{ab} \right|^2 
\] (16)

and given by the integral over all times \( t > 0 \) retarded flow of the particles reemitted by the environment

\[
\sigma^{ab}_2(E) = \Gamma_s \int_0^\infty dt \left| \left( \frac{A^\dagger}{E - \mathcal{H}^{(s)} + \frac{i}{2}\Gamma_s} e^{-i\mathcal{H}^{(s)} t} A \right)^{ab} \right|^2. 
\] (17)

The contribution (16) is identical to the result of the Efetov’s imaginary-potential model if we identify the strength of this potential as \( \alpha = \frac{\pi}{2} \Gamma_s \). The second contribution (17) compensates the loss of the flow due to absorption described by the first one.
For higher scattering energies, one should take into account that excited background states are not stable and can strongly overlap, $\gamma_e \gg d$, even with no coupling to the doorway states. This smears out the fine-structure fluctuations thus reducing the loop functions to $g^{(e)}(E) \Rightarrow -i$ and $l^{(e)}(E) \Rightarrow 2/\gamma_e$ [15]. As opposed to the above consideration, the unitarity of the scattering submatrix $S$ is broken. The structures narrower than $\gamma_e$ are excluded and all particles delayed for a time larger than $1/\gamma_e$ are irreversibly absorbed and lost from the outgoing flow.

4. It follows that there exist two different temporal characteristics of the scattering process [2]. The first one is the decay rate $\gamma_0 + \Gamma_s$ of the doorway state once excited through the incoming channel. This rate is readily seen from the scattering amplitude $\mathcal{T}(E) = \gamma_0 \left[ E - \varepsilon_0 + \frac{i}{2} (\gamma_0 + \Gamma_s) \right]^{-1}$ near an isolated doorway resonance. Since the doorway state is not an eigenstate of the total effective Hamiltonian $\mathcal{H}$, it fades exponentially out not only because the particle is emitted onto the outer channels but also due to the internal transitions with formation of the exact fine structure resonances over which the doorway state spreads. More than that, even if the background is stable the internal decay remains exponential up to the Heisenberg time $2\pi/d$. Typically, during the time $\tau_s = 1/\Gamma_s$ the background absorbs the particle. After this, it can evade via one of the $e$-channels or be after a while emitted back into the dot and finally escape onto an outer channel. Since the particle reemitted by the chaotic background carries no phase memory, the characteristic time during which this memory is lost (dephasing time) coincides with $\tau_s$. All the resonances, save the broad one with the width $\Gamma_0 = \gamma_0 - \Gamma_s$, have rather small widths and decay much slower. Just the resonances with the widths within the interval $\gamma_e < \Gamma_a \ll \Gamma_s$ contributes principally in the Wigner time delay which shows how long the excited system still returns particles into the observed channel. For example the delay time near the energy of an isolated doorway resonance equals

$$\tau_W(E) = \frac{\gamma_0}{E^2 + \frac{1}{4} (\gamma_0 + \Gamma_s)^2 \Lambda}. \quad (18)$$

This result differs by the enhance factor $\Lambda = 1 + \Gamma_s/\gamma_e$ (compare with eq. [13]) from what follows from the somewhat oversimplified consideration in [16], which implies that interaction with an environment yields only absorption.

Turning to general consideration, we model as usual the unperturbed chaotic single-particle motion using the random matrix theory. The observable channels are considered
below to be statistically equivalent and, to simplify the calculation, no time-reversal symmetry is suggested. We suppose below that $d \ll \gamma_e \ll \min(D, \Gamma_s) \ll 1$ (the radius of the semicircle). If, on the contrary, $\gamma_e \gg \Gamma_s$ the transitions into the background become equivalent to irreversible decay similar by its properties to the decay into continuum. In such a limit of full absorption the factor $\Lambda \to 1$ and the approach of ref. [16] becomes valid.

It is readily seen that to account for the interaction with the background the substitution $\varepsilon \Rightarrow \varepsilon - i\Gamma_s$ should be done while calculating the two-point S-matrix correlation function $C(\varepsilon) = S(E) \otimes S^\dagger(E+\varepsilon)$. This immediately yields the connection $C_V(t) = \exp(-\Gamma_st)C_0(t)$ between the Fourier transforms with and without interaction, the additional damping being as before caused by the spreading over the fine-structure resonances.

The modification given in [16] of the method proposed in [17] allows us to calculate also the distribution $P(q) = (\pi M^{(s)})^{-1}\text{Im}\langle \text{tr}(q - \mathbf{Q} - i\mathbf{0})^{-1} \rangle$ of proper delay times (eigenvalues of the Smith matrix). The corresponding generating function is proportional to the ratio of the determinants of two $2N^{(s)} \times 2N^{(s)}$ matrices with the following structure (compare with [16])

$$A(z) = -i \left( E - H^{(s)} \right) \sigma_3 + \frac{1}{2} \left( AA^\dagger + \Gamma_s \right) - \frac{\Lambda}{z} \left( 1 + \sqrt{1 - \frac{\Gamma_s}{\Lambda} z \sigma_1} \right) \approx -i \left( E - H^{(s)} \right) \sigma_3 + \frac{1}{2} AA^\dagger - (1 + \sigma_1) \left( \frac{\Lambda}{z} - \frac{\Gamma_s}{2} \right)$$

where the variable $z$ spans the complex $q$-plane. The square root sets [16] the restriction $q \leq \Lambda/\Gamma_s \approx 1/\gamma_e$ on the positive real axes. In the approximation of the second line valid if $\Lambda \gg 1$ we arrive to a simple relation

$$P_V(\tau) = \frac{1}{\Lambda} P_0(\tau_V), \quad (20)$$

where $\tau_V \equiv (\tau/\Lambda) [1 - \pi (\Gamma_s/\Lambda D) \tau]^{-1} > 0$ and $\tau = \frac{D}{2\pi q}$. This relation does not hold in the asymptotic region $\tau_V \to \infty$ or $\tau \to \Lambda D/\pi \Gamma_s \approx \frac{1}{\pi} D/\gamma_e$ where the more elaborate rigorous expressions obtained in [16] must be used with the substitution $\tau \Rightarrow \tau/\Lambda$ being made. For example in the single-channel scattering with the perfect coupling to the continuum the time delay distribution (see [19]) $P_V(\tau) = e^{-1/\tau_V} / \tau_V^3$ reaches its maximum at the point $\tau_V = 1/3$ or $\tau = (1/3) \Lambda \left( 1 + \frac{2}{3} \Gamma_s/D \right)^{-1} \gg 1/3$. The most probable delays shift towards larger values. The estimation just made is quantitatively valid only if $\Gamma_s \ll D/2\pi$ and the maximum lies near the point $\frac{1}{3} \Gamma_s/\gamma_e$ distant from the exact edge $D/2\pi\gamma_e$ of the distribution. Under the
latter restriction the approximate formula (20) describes well the bulk of the delay time distribution which becomes, roughly, \( \Lambda \) times wider. The condition noted is much less restrictive in the case of weak coupling to the continuum when the transmission coefficient \( T \ll 1 \) and the most probable delay time \( \tau \approx T \Gamma_s/4\gamma_e \ll D/2\pi\gamma_e \) as long as \( \Gamma_s \ll D/T\pi \).

The approximation (20) works even better when the number of channels \( M^{(s)} \gg 1 \) and the delay times are restricted to a finite interval \([17, 18]\). The delay time scales in this case with \( M^{(s)}^{-1} \) and the natural variable looks as \( \tau = qM^{(s)}D/2\pi = \Gamma_Wq/T \) with \( \Gamma_W \) being the Weisskopf width. The edges of the distribution (20) are displace towards larger delays by the factor \( \Lambda \), \( \tau_{V(\pm)} = \Lambda\tau_{0(\pm)} \). One can readily convince oneself that the taken approximation remains valid in the whole interval \( \Delta\tau_V = \Lambda(\tau_{0(+) - \tau_{0(-)})} \) under condition \( \Gamma_s \ll \Gamma_W \) which implies weak interaction with the background. The width of the delay time distribution broadens by the factor \( \Lambda \) due to the influence of the background.

5. In summary, the influence of a complicated environment on the chaotic single-particle scattering is analyzed. Unlike some earlier considerations, the coupling to the background is supposed to be purely Hermitian. The single-particle doorway states which are excited inside a mesoscopic device and observed through external channels are additionally damped with the rate \( \Gamma_s = 2\pi\langle |V|^2 \rangle/d \) because of the spreading over the long-lived fine-structure resonances. This rate uniquely determines the dephasing time during the particle transport through the ballistic microstructure. Absorption takes place because of hidden decays of the background resonance states. The formation of the fine-structure resonances strongly enhances delay time fluctuations thus, in particular, broadening the distribution of the proper delay times.

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