Contents

DIMENSIONAL REDUCTIONS OF GRAVITY AND RELATIONS BETWEEN STATIC STATES, COSMOLOGIES AND WAVES

A. T. Filippov

Various approaches to dimensional reduction of gravity theories to two-dimensional and one-dimensional dilaton gravity models were considered in [1]-[3]. In particular, an unusual dimensional reduction of an integrable (1+1)-dimensional dilaton gravity coupled to matter was proposed. The reduction works in the moduli space of the (1+1)-dimensional theory reducing the general space-time dependent solutions to one-dimensional static states (black holes in particular), cosmological models, and waves. An unusual feature of this reduction is the fact that the wave solutions depend on two variables – space and time. It was shown that these waves can also be derived by a generalized separation of variables (applicable also to non-integrable models and to higher dimensional theories). These results clearly demonstrate a nontrivial and close relation between static states, cosmologies, and waves generalizing earlier proposed static - cosmological duality.

Using the generalized separation of variables the one-dimensional states depending on both the space and time variables were found in the spherically symmetric higher-dimensional gravity coupled to scalar matter fields. This procedure is more general than the usual 'naive' reduction and apparently more general than the reductions that employ group theoretical methods. In this way, unusual generalizations of the spherical static states and cosmologies were constructed that are related by a simple duality transformation. Even more interesting new static states, cosmologies and waves possibly exist in cylindrical and axial theories. A generalization of the standard cylindrical four-dimensional gravity was proposed, which reduces to a (1+1)-dimensional dilaton gravity with nontrivial cosmological potentials produced by a four-dimensional Kaluza-Klein mechanism. Such potentials do not exist in the spherical theory but may be present in any axially symmetric theory even if it is almost spherical. This means that any axial, almost spherical universe is qualitatively different from the exactly spherical one. In particular, in the almost spherical universe, geometry may generate weak cosmological potentials that possibly could imitate dark energy effects.

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QUANTUM-GROUP (SUPER)SYMMETRIES IN NONCOMMUTATIVE (SUPER)FIELD THEORIES

B.M. Zupnik

Noncommutativity of the space-time coordinates naturally emerges in the low-energy limit of the string theory on the special backgrounds. Field theories on the simplest noncommutative space-time can be described via the Moyal-Weyl noncommutative $\star$-product defined on the tensor product of ordinary commutative functions

$$f \star g = \mu \circ \exp(\frac{i}{2}\vartheta^{mn}\partial_m \otimes \partial_n) f \otimes g = \lim_{x=y} \exp(\frac{i}{2}\vartheta^{mn}\frac{\partial}{\partial x^m}\frac{\partial}{\partial y^n}) f(x)g(y),$$

where ϑ^{mn} are constant deformation parameters. The noncommutative (NC) field theory uses $\star$-products of scalar, spinor or vector fields in each term of the action, and this selection rule does not follow from the standard Lorentz symmetry which is broken in the noncommutative interactions.

It was recently shown [1] that this NC field theory preserves a new quantum-group type symmetry (twisted Poincaré group) which uses the twist operator $\mathcal{F}_P$ in the co-product Δ_t acting on the tensor product of fields

$$\mathcal{F}_P = \exp(\frac{i}{2}\vartheta^{mn}\partial_m \otimes \partial_n), \quad \Delta_t(M_{mn}) = \mathcal{F}_P(M_{mn} \otimes 1 + 1 \otimes M_{mn})\mathcal{F}_P^{-1},$$

where M_{mn} are the Lorentz group generators. The co-product Δ_t generates the covariant twisted transformation of the $\star$-product

$$\hat{\delta}_\omega(f \star g) = \frac{1}{2}\omega^{mn} M_{mn}(f \star g).$$

All actions using this $\star$-product are invariant under the twisted Poincaré transformations, and the corresponding free actions preserve also the standard Poincaré symmetry.

The same operator $\mathcal{F}_P$ can be used to deform the general covariance (diffeomorphism) group of the gravity theory [1, 2]

$$\Delta_t(\xi) = \mathcal{F}_P(\xi \otimes 1 + 1 \otimes \xi)\mathcal{F}_P^{-1}, \quad \xi = \xi^m(x)\partial_m,$$

where $\xi^m(x)$ are arbitrary functions. Note that $\star$-products of scalar, vector or tensor fields transform covariantly in the deformed diffeomorphism quantum group, while the standard Leibniz rules are not satisfied for the twisted variations of these products. The deformed noncommutative gravity theory was constructed in this approach.

We proposed a new version of the NC gravity formalism based on the same twisted diffeomorphism group $\hat{\delta}_\xi(f \star g) = -\xi(f \star g)$ [3]. In contrast to Refs.[1, 2], our method preserves reality at each stage of calculations in the NC gravity. For instance, we use the real representation for the metric $g_{mn} = \eta_{mn} + \kappa h_{mn}$ (κ is the gravity coupling constant, η_{mn} is the Minkowski space metric) and construct real NC generalizations of the Christoffel symbols and Riemann tensor as $\star$-polynomials in κh_{mn}. These constructions are covariant under the background twisted Poincaré symmetry. The NC Riemann tensor satisfies the generalized Bianchi identities. Our minimal representation of the NC gravity action and NC generalization of the Einstein equations contain only the ϑ^{mn}-corrections of the even degrees in comparison with the standard gravity action and equations. Of course, one can add the independent pure noncommutative action terms of the odd degree in the deformation parameters.

The simplest non-anticommutative (NAC) supersymmetric field models were studied in the framework of the deformed Euclidean superspaces (see, e.g. the review [4]). The Q-deformation of the supercommutative product of superfields A and B can be defined via the twist operator $\mathcal{F}_Q$ acting on the tensor product of the $N=(1,1)$ superspaces [5]

$$A\star B=\mu\circ\mathcal{F}_Q(A\otimes B),\quad \mathcal{F}_Q=\exp(-\frac{1}{2}C^{\alpha\beta}_{ik}Q^i_\alpha\otimes Q^k_\beta),$$

where Q^i_α are the left-handed spinor and the SU(2) isospinor generators of the Euclidean $N=(1,1)$ supersymmetry and $C^{\alpha\beta}_{ik}$ are the deformation constants. This $\star$-product of superfields is not covariant with respect to the standard $\mathrm{SU}_R(2)$rotations of the right-handed spinors and the corresponding right supersymmetries $\bar{Q}_{k\dot\alpha}$. The Q-deformation preserves chiral, antichiral and Grassmann analytic representations of the $N=(1,1)$ supersymmetry. The singlet Q-deformation corresponds to the choice $C^{\alpha\beta}_{ik}=2I\varepsilon^{\alpha\beta}\varepsilon_{ik}$, it is invariant under the $\mathrm{SU}_L(2)\times\mathrm{SU}_R(2)\times\mathrm{SU}(2)$ group. Supersymmetric models with the singlet Q-deformation were analyzed in [6]-[8]. In particular, we proved the ultraviolet finiteness of the Q-deformed $N=(1,1)$ electrodynamics with the specific noncommutative interaction, induced by the singlet deformation constant I.

The quantum-group structure of the deformed $N=(1,1)$ supersymmetry is defined by the operator $\mathcal{F}_Q$ [5]. It is possible to choose the standard representation for all generators of the $N=(1,1)$ supersymmetry acting on the primary superfields A,B. The twisted supersymmetry acts covariantly on the $\star$-product of superfields

$$\hat\delta_{\bar\epsilon}(A\star B)=\bar\epsilon^{k\dot\alpha}\bar{Q}_{k\dot\alpha}(A\star B)\neq(\hat\delta_{\bar\epsilon}A\star B)+(A\star\hat\delta_{\bar\epsilon}B).$$

Thus, the actions of deformed superfield theories [6]-[8] using the $\star$-product of superfields are invariant with respect to the transformations of the twist-deformed $N=(1,1)$ supersymmetry. The quadratic free superfield actions are also invariant under the standard $N=(1,1)$ supersymmetry.

The twist-deformed symmetries and supersymmetries provide the selection rules for the deformed noncommutative (super)field theories. We hope that these quantum-group type symmetries will be useful in constructing realistic noncommutative models of gravity and supergravity.

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HIGHER SPINS IN EXTENDED SPACES

S.A. Fedoruk, E.A. Ivanov

Nowadays, there is a growing evidence that the higher spin (HS) field theory (see [1, 2] and refs. therein) can play a central role in the continuing quest for the Fundamental Theory. The HS theory deals with infinite sets of massless fields of all spins and possesses some generalized (super)conformal symmetry. It was conjectured that spontaneous breaking of the underlying (super)symmetries of this fundamental system, possibly accomplished in several successive steps, could yield the string–like phase with massive higher spins coupled to (super)gravity. In other words, String Theory might happen to be none other than a spontaneous breaking phase of more fundamental HS theory. There are many indications that an adequate geometric arena for treating HS fields and their symmetries is provided by extended (super)spaces with additional bosonic twistor-like and/or tensorial coordinates.

One of the important symmetries realized on $4D$ massless HS fields is a generalization of the standard conformal symmetry $SU(2,2)$. In the bosonic limit it is the $Sp(8)$ symmetry, and it extends to the supergroup $OSp(1|8)$ in the full supersymmetric variant of HS theory.

In [3], surprising relations were exhibited between the HS theory and nonlinear realizations of the supergroup $OSp(1|8)$ in the supercoset $OSp(1|8)/SL(4,R)$. The latter involves the tensorial superspace, parametrized by usual $N=1$, $4D$ superspace coordinates x^μ, θ^α, $\bar\theta^{\dot\alpha}$ and six commuting tensorial coordinates $z^{\alpha\beta}$, $\bar z^{\dot\alpha\dot\beta}$, as well as some Goldstone superfields. Imposing some covariant constraints allows one to algebraically eliminate all Goldstone superfields in terms of a single superfield $\Phi(x,z,\bar z,\theta,\bar\theta)$ associated with the dilatation generator. Simultaneously, these constraints imply the dynamical equations for the basic superfield

$$D^\alpha D_\alpha\,\Phi = \bar D^{\dot\alpha}\bar D_{\dot\alpha}\,\Phi = 0\,, \qquad [D_\alpha, \bar D_{\dot\alpha}]\,\Phi = 0\,, \tag{1}$$

which coincide with the set derived in [4] while studying wave equations of tensorial superparticle [5]. Thus one novel point of the nonlinear realization approach [3] is that the basic superfield encompassing all spins appears as a parameter of the supercoset $OSp(1|8)/SL(4,R)$. Another point is the new geometric interpretation of the associate HS equations. They proved to be the conditions of vanishing of the covariant spinor derivatives of the Goldstone superfields related with the generators of dilatations and conformal supersymmetry.

There exists another way to generalize (super)conformal group to higher spins. It consists in extending $SU(2,2)$ to $SU(3,2)$ and supersymmetrizing the latter to the supergroup $SU(3,2|1)$.

The model of HS particle possessing $SU(3,2)$ symmetry was proposed in [6]. This particle moves in Minkowski space extended by commuting Weyl spinor ζ^α, $\bar\zeta^{\dot\alpha}$. The part of the $SU(3,2)$ group symmetry which lies out of $SU(2,2)$, is realized as the even counterpart of $N=1$, $4D$ supersymmetry translations

$$\delta x^{\dot\alpha\alpha} = i(\bar\epsilon^{\dot\alpha}\zeta^\alpha - \bar\zeta^{\dot\alpha}\epsilon^\alpha)\,, \quad \delta\zeta^\alpha = \epsilon^\alpha\,, \quad \delta\bar\zeta^{\dot\alpha} = \bar\epsilon^{\dot\alpha} \tag{2}$$

(with the *commuting* transformation parameter ϵ^α) and as the even counterpart of superconformal boosts. The spectrum of the particle of ref. [6] contains all helicities, every non-zero helicity appearing only once.

In [7] this particle model was extended to a superparticle possessing $SU(3,2|1)$ supersymmetry, a closure of the standard (odd) four-dimensional $N=1$ superconformal symmetry $SU(2,2|1)$ and its even $SU(3,2)$ counterpart. This generalization allows one to preserve the important notion of $N=1$ chirality whereas in the $OSp(1|8)$ case the chirality can be introduced only at cost of adding some extra harmonic variables [3]. Specifically, the space of states of the model of [7] is spanned by an infinite set of massless supermultiplets combined into a single chiral superwave function $\Psi(\hat{x}_{L},\zeta,\theta)$, $\hat{x}_{L}^{\dot\alpha\alpha}=x^{\dot\alpha\alpha}+i\bar\zeta^{\dot\alpha}\zeta^{\alpha}+i\bar\theta^{\dot\alpha}\theta^{\alpha}$. It satisfies the superconformally covariant equations

$$D^{\alpha}D_{\alpha}\,\Psi=0\,,\qquad D^{\alpha}\nabla_{\alpha}\,\Psi=0\,,\tag{3}$$

where ∇_α is the even counterpart of the $N=1,4D$ spinor covariant derivative D_α. Eqs. (3) encompass both the dynamical equations and Bianchi identities for the massless chiral superfield strengths with arbitrary external helicities in the ζ expansion of the HS superfield $\Psi(\hat{x}_{L},\zeta,\theta)$.

The "master" HS particle model was proposed in [8]. It makes manifest a classical equivalence of the HS particle corresponding to the unfolded formulation of massless HS fields [1] and the HS particle model with the even counterpart of supersymmetry [6, 7]. The master HS particle propagates in a space which is parametrized, in addition to x^μ, ζ^α, by twistor-like spinor variables λ_α, y^α and an extra complex scalar η. In the variables $x_{L}^{\dot\alpha\alpha}=x^{\dot\alpha\alpha}+i\,\bar\zeta^{\dot\alpha}\zeta^{\alpha}$, $\bar y_{L}^{\dot\alpha}=\bar y^{\dot\alpha}+2i\,\eta\,\bar\zeta^{\dot\alpha}$, the quantum counterparts of the classical spinor constraints express all fields in the expansion of the wave function over ζ^α, $\bar\zeta^{\dot\alpha}$ in terms of single field $\Phi^{(q)}(x_L,y,\bar y_L,\eta)$ obeying the modified unfolded equation

$$\left(\partial_{L\,\alpha\dot\alpha}+i\frac{\partial}{\partial y^{\alpha}}\frac{\partial}{\partial\bar y_{L}^{\dot\alpha}}\right)\Phi^{(q)}=0\,.\tag{4}$$

The new scalar $U(1)$ constraint implies the wave function to possess an external $U(1)$ charge q defined as a degree of homogeneity with respect to both spinor and scalar coordinates. In the holomorphic expansion w.r.t. the scalar coordinate

$$\Phi^{(q)}(x_L,y,\bar y_L,\eta)=\sum_{k=0}^{\infty}\eta^{k}\phi^{(q+k)}(x_L,y,\bar y_L)\,,\tag{5}$$

due to eq. (4), every field $\phi^{(q+k)}(x_L,y,\bar y_L)$ includes only one independent complex self–dual field strength $\phi_{\alpha_1\ldots\alpha_{q+k}}(x_L)$ of the massless particle of helicity $\frac{q+k}{2}$. Thus, depending on their external $U(1)$ charge q, the higher spin fields $\Phi^{(q)}$ accommodate different higher spin multiplets. The all-helicity higher spin multiplet of the unfolded formulation [1] (with a complex scalar field) is recovered at $q=0$. At $q\neq 0$ there emerge new higher spin multiplets having different helicity contents. For $q>0$ they are constituted by self–dual field strengths of growing positive helicities, starting from $q/2$. When $q<0$, the spectrum contains an infinite tower of massless states of all positive helicities and, in addition, a finite number of states with negative helicities starting from $q/2$ ('spin–flip' multiplets). In [8], an equivalent description of these multiplets in terms of chiral "superfields" of the even supersymmetry was also given.

A further exploration of HS fields along these promising directions is now in progress.

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UNIVERSAL BETHE ANSATZ

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Quantum inverse scattering method or algebraic Bethe ansatz was designed to address the eigenvalue problem for the mutually commuting integrals of motion in a quantum integrable system, which possesses the hidden symmetry. Up to now this is the most powerful method to describe the space of states in the integrable models with non-trivial interactions.

This space of states is characterized by the properties of the concrete model as well as by the symmetry of the system. It is a challenge problem to find the restrictions for this space in the class of integrable models which share the same or similar symmetry.

The fundamental object for any classical or quantum integrable model is so called L-operator which depends on the local dynamical variables of the system. Taking the ordered product or integral of these operators one may construct the monodromy matrix which is a source of the integrals of motion. The commutation relations for the dynamical variables are translated into commutation relations for the elements of the monodromy matrix which has the standard form of $RTT = TTR$ relations in the case of quantum integrable models. Symmetry of the model is characterized by the type of R-matrix entering into this commutation relation and algebraic Bethe ansatz allows one to describe the vectors of the space of states as certain combinations of the matrix elements of monodromy matrix.

Solution of this problem is divided into two steps. First, one has to describe the structure of the eigenvectors and, second, to impose the so-called Bethe equations for the parameters of Bethe vectors. The vectors with free parameters are called off-shell Bethe vectors. They are an important ingredient in the theory of deformed or quantized Knizhnik-Zamolodchikov equations.

The structure of the off-shell Bethe vectors in the case of integrable models with symmetry of minimal rank is trivial. In this case, the monodromy matrix is two by two matrix and off-shell Bethe vectors are the products of the off-diagonal elements. In the case of higher rank symmetries, the structure of the off-shell Bethe vectors are defined by the hierarchical Bethe ansatz. This method gives the description of the off-shell Bethe vectors in the implicit form, namely, the problem for the symmetry algebra of the rank r is reduced to an analogous problem for the rank $r-1$. The hierarchical Bethe ansatz yields quite easily the Bethe equations, but does not produce effective formulae for the off-shell Bethe vectors.

A theory for the construction of the off-shell Bethe vectors in quantum integrable models with arbitrary symmetry was developed in [1, 2, 3, 4]. To formulate this theory we used the well-known fact that the $R(z_1/z_2)T_1(z_1)T_2(z_2) = T_2(z_2)T_1(z_1)R(z_1/z_2)$ commutation relation for a given R-matrix defines a certain infinite dimensional algebra, which is generated by the modes of the matrix elements of the universal monodromy matrix $T(z)$. This algebra has a direct relation to the symmetries of the concrete model which corresponds to certain representations of this algebra. In [2], the most general properties of the off-shell Bethe vectors were reformulated in terms of projections onto intersections of Borel subalgebras of the algebra generated by the universal monodromy matrix. In [1], the methods for calculation of these projections for the class of quantum affine algebras were developed. In [3, 4], the theory was compared with known constructions of off-shell Bethe vectors in the case of the quantum affine algebra $U_q(\widehat{gl}_N)$.

The results obtained in these papers allow one to combine hierarchical and analytical Bethe ansatz into a unite theory which was called the universal Bethe ansatz method.

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ENTANGLEMENT, HOLOGRAPHY, AND QUANTUM GRAVITY

D.V. Fursaev

Quantum entanglement is an important physical phenomenon in which the quantum states of several objects cannot be described independently, even if the objects are spatially separated. Quantum entanglement is used in different research areas [1]. In quantum information theory entangled states are a valuable source of processing the information. Quantum entanglement also plays an important role in properties of strongly correlated many-body systems and in collective phenomena such as quantum phase transitions.

In [2] a general geometrical structure of the entanglement entropy for spatial partition of a relativistic QFT system was established by using methods of the effective gravity action and the spectral geometry. A special attention was payed in [2] to the subleading terms in the entropy in different dimensions and to behaviour in different states.

The work [2] has also dealt with another problem related to the entanglement in a fundamental (quantum gravity) theory. It was conjectured that for partition of the ground state by a plane the entanglement entropy per unit area of the plane is

$$\mathcal{S} = \frac{1}{4G_N} \quad , \tag{1}$$

where G_N is the gravitational constant in the low-energy gravity sector of the theory. $\mathcal{S}$ measures entanglement of all genuine microscopical degrees of freedom of the fundamental theory.

If the fundamental theory lives in a space-time with number of dimensions higher than four there is an issue of higher-dimensional extension of (1). There may be at least two options. In theories of the Kaluza-Klein (KK) type all degrees of freedom propagate in the bulk. In this case partition of all KK modes by a hypersurface $\mathcal{B}$ in four-dimensions is equivalent to partition of higher-dimensional fields by the surface $\tilde{\mathcal{B}} = \mathcal{B} \times \Omega_{D-4}$ where Ω_{D-4} is a small compact space of $D-4$ extra-dimensions. If (1) is true the entanglement entropy should be

$$S = \frac{\mathcal{A}}{4G_N} = \frac{\tilde{\mathcal{A}}}{4G_N^{(D)}} \quad , \tag{2}$$

where $\mathcal{A}$ is the area of $\mathcal{B}$ and $\tilde{\mathcal{A}} = \mathcal{A} \cdot \mathrm{vol}\ \Omega_{D-4}$ is the volume of $\tilde{\mathcal{B}}$. Equation (2) takes into account the relation $G_N^{(D)} = G_N \cdot \mathrm{vol}\ \Omega_{D-4}$, between higher-dimensional, $G_N^{(D)}$, and low-dimensional, G_N, gravitational couplings.

The other option is the theories with large extra dimensions. In these theories the gravity remains higher-dimensional while matter fields are confined on a brane and do not propagate in the bulk. For this reason, the way how the separating surface is extended in extra-dimensions should be determined by the bulk gravity equations.

In the Randall-Sundrum model the bulk is the anti-de Sitter (AdS) gravity. If $\mathcal{B}$ is a separating hypersurface in the brane theory, its extension $\tilde{\mathcal{B}}$ has to be a surface embedded in the AdS space. The boundary of $\tilde{\mathcal{B}}$ on the brane is $\mathcal{B}$. With this prescription the entropy in the brane theory takes the form (2), where $G_N^{(D)} = G_N^{(d+1)}$ is the gravitational constant in the bulk and d is the number of dimensions on the brane.

Relation $S = \tilde{\mathcal{A}}/(4G_N^{(d+1)})$ in this setting can be considered as a "holographic formula". It was first suggested by Ryu and Takayanagi [3], [4] for the entropy of conformal field theories (CFT) which admit a dual description in terms of the AdS gravity one dimension higher. The authors showed that $\tilde{\mathcal{B}}$ has to be a *minimal* surface in the bulk and presented an intuitive derivation of the formula based on the AdS/CFT correspondence.

A general proof of the holographic formula was given in [5]. The entanglement entropy was determined by a partition function which was defined as a path integral over Riemannian AdS geometries with non-trivial boundary conditions (the topology of the Riemannian spaces puts restrictions on the choice of the minimal hypersurface for a given boundary conditions).

The fact that quantum entanglement and gravity are connected phenomena is interesting for several reasons. First, on the base of (1) one can study properties of the gravitational couplings in "gravity analogs by using entanglement entropy in condensed matter systems. Second, the holographic formula $S = \tilde{\mathcal{A}}/(4G_N^{(d+1)})$ enables one to study quantum entanglement in strongly correlated systems, such as strongly coupled gauge theories, where the direct field theoretical computations for the entropy are extremely difficult. This can be done by purely geometrical methods by studying minimal surfaces in anti-de Sitter spaces.

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QUASI-NORMAL MODES OUTSIDE THE BLACK HOLE PHYSICS

V.V. Nesterenko

Quasi-normal modes (QNM) are widely used now in black hole physics and in relativistic theory of stellar structure [1]). The corresponding eigenfrequencies are complex numbers, however it is not due to the dissipative processes but it is a consequence of unbounded region occupied by the oscillating system. The latter naturally leads to the energy loss due to the wave emission (for example, gravity waves).

However QNM are not the peculiarities of the gravitational problems only. Actually, they appear, in a natural way, when considering the oscillating systems unbounded in space (Gamov states in quantum mechanics, quasi-bound states in the theory of open electromagnetic resonators). The necessary condition for emergence of such modes is imposing the *radiation condition* at spatial infinity on the field functions. It is this condition that leads to the characteristic behaviour of the quasi-normal modes, namely, these solutions to the relevant equations (usually Helmholtz equation) exponentially decay in time when $t \to \infty$ and simultaneously they exponentially rise at spatial infinity $r \to \infty$. By the way, the radiation of a portable telephone excites the electromagnetic QNM inside a human body, first of all, inside the human head, This fact should be taken into account when estimating the pertinent health hazard [3].

In papers [2, 3] it was proposed to investigate the spatial form of the quasi-normal modes with allowance for their time dependence. Indeed, these solutions have the character of propagating waves that are eventually going to spatial infinity. Let us take the point which is sufficiently far from the region with nontrivial dynamics in the system under study. Obviously, it has sense to say about the value of the quasi-normal mode at a given point only after arrival at this point of the wave described by this mode. The maximal value of the quasi-normal mode is observed just at the moment of its arrival at this point. At the later moments the quasi-normal mode is dumping due to its characteristic time dependence. With allowance of all this the maximal observed value of the quasi-normal mode proves to be finite. Thus, the exponential rising of QNM at infinity is *not observable*.

When we have QNM instead the usual normal modes it implies that we are dealing with the *open systems* when the energy can be radiated to infinity. Open systems admit a dual description: on the one hand, they can be considered from the “inside” point of view, treating the coupling to the environment as a – not necessarily small – perturbation. From this point of view, one can study the (discreet) complex eigenvalues of the system, the width of resonances and the resulting decay properties. On the other hand, open systems allow to take the “outside” point of view, considering the system as a perturbation of the environment. The typical quantity to be investigated in this approach is the *scattering matrix* (S-matrix), i.e. the amplitude for passing from a given incoming field configuration to a certain outgoing configuration as a function of energy which is strictly positive but continuous. The equivalence of these two descriptions has been proved in a recent paper [2].

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