Beam-like Excitations of Kerr-Schild Geometry and Semiclassical Mechanism of Black-Hole Evaporation.

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It has been observed that exact solutions for electromagnetic (EM) excitations of the Kerr-Schild (KS) geometry form outgoing beams which have very strong back reaction to metric and break the black hole horizon. As a result, interaction of a black hole with nearby electromagnetic field and electromagnetic vacuum has to cover the horizon by a set of fluctuating microholes. We integrate and analyze the Debney-Kerr-Schild equations for electromagnetic excitations of a black-hole and obtain that the exact solutions for outgoing radiation contain two related but radically different components which shed light on a possible semi-classical mechanism of black-hole evaporation: a) first component consists of the singular beam pulses which perforate horizon, breaking its impenetrability, and b) another component is regular and responsible for the loss of mass similar to the known Vaidya ‘shining star’ radiation. We show also that the mysterious twosheeted twistor structure of the Kerr-Schild geometry corresponds to a holographic structure of quantum black hole spacetimes predicted by Stephens, t’ Hooft and Whiting. The resulting Kerr-Schild geometry of fluctuating twistor-beams takes an intermediate position between the classical and quantum gravity.

I. INTRODUCTION

Diversity of the recent ideas on the origin of black hole evaporation [1] has a common core based on complex analyticity and conformal field theory, which unifies the black hole physics with (super)string theory and physics of elementary particles, as it was first noticed by ‘t Hooft in [2]. These relationships open a way to quantum gravity by a holographic correspondence between a classical bulk gravity and a quantum Conformal Field Theory which lives on a holographically dual to bulk 2D boundary [3, 4, 5, 6].

In this paper we show that analytic twosheeted structure of the Kerr-Schild geometry is in perfect agreement with the required holographic structure of quantum black hole spacetimes [7, 8]. Obtaining exact solutions of Debney-Kerr-Schild equations [9] for electromagnetic excitations of the Kerr-Schild geometry and their back reaction to metric and horizon, we arrive at a semiclassical fluctuating Kerr-Schild geometry which takes an intermediate position between the classical and quantum gravity. This geometry has a classical fluctuating fine-grained structure which shows that the usual representations on the structure of BH and the horizon may be naive and very far from reality [30].

The usual statements on stability of the black hole horizon are based on the theorems (Robinson and Carter) [11] claiming the uniqueness of the Kerr solution. These theorems are valid under a series of conditions: stationarity, axial symmetry and asymptotic flatness of the metric which turn out to be broken under electromagnetic (EM) excitations of black hole [12]. Another vulnerable point is the traditional use of perturbative approach. The recent analysis of the exact nonstationary electromagnetic solutions, performed in the Kerr-Schild formalism [13, 14], showed that excitations of black hole do not contain the smooth spherical harmonics used by the perturbative analysis: elementary exact electromagnetic excitations on the Kerr background have the form of singular beams which may be outgoing in any angular direction which have very strong back reaction on metric and the horizon. The beams break horizon topologically [14, 15], forming the holes which allow matter to escape interior of black hole. Origin of this effect lies in analyticity of the Kerr-Schild solutions and, in particular, in twistor analyticity of the Kerr congruence determined by the Kerr theorem [16, 17, 18]. The Kerr-Schild form of metric

\[ g_{\mu\nu} = \eta_{\mu\nu} + 2H_k\mu\nu, \]

has many advantages with respect to other representations. First of them is the rigid connection of coordinates \( x^\mu \) to auxiliary Minkowski space-time with metric \( \eta_{\mu\nu} \) and the corresponding unfastening of the coordinates and solutions from positions of the horizons, which allows one to analyze deformations of the horizon. Second, the resulting Kerr-Schild solutions are not singular at the horizon. Third advantage is related with the Kerr-Schild twistor structure [18] and correspondence with the requirements of holographic principle and the presumable properties of a quantum black hole space-time described by Stephens, t’ Hooft and Whiting in [8]. The Kerr-Schild holographic projection is realized by the Kerr congruence and generalizes the holographic map considered by Bouso [19]. The twosheeted Kerr-Schild structure perfectly matches with holographic approach and allows one to consider evaporation as a scattering of in-vacuum on a holographically dual 2+1 dimensional Kerr source which separates the in and out regions [8], and we obtain that the exact outgoing solutions have two related but radically different components which shed light on a possible semi-classical mechanism of black-hole evaporation:

a) a casual set of singular beam pulses which perforate horizon, breaking its impenetrability, and

b) a regular component which is responsible for the loss of mass and akin to the known Vaidya ‘shining star’ radiation.
For the reader convenience we give here some digression with description of the Kerr-Schild relations [9]. Analyticity of the Kerr-Schild geometry originates from the complex function $Y = e^{i\phi} \tan \frac{\theta}{2}$ which is a projection of celestial sphere $S^2$ on the complex plane. This function determines the Kerr congruence and complex tetrad forms. The Kerr theorem sets the dependence $Y(x)$, and the null vector field $k^\mu$ of the Kerr-Schild metric form is expressed via function $Y(x)$

$$k_\mu dx^\mu = P^{-1}(du + \bar{Y} d\zeta + Y d\bar{\zeta} - Y \bar{Y} dv), \quad (2)$$

in the null Cartesian coordinates $2^\zeta = x + iy$, $2^\bar{\zeta} = x - iy$, $2^u = z - t$, $2^v = z + t$. Therefore, the field $k^\mu(x)$, $x \in M^4$ determines symmetry of space-time, polarization of the Kerr-Newman electromagnetic field, direction of gravitational ‘dragging’ and so on. This vector field is tangent to the Kerr congruence which is the family of the light-like geodesic lines, in fact twistors. Twisting structure of the Kerr congruence is shown in Fig.1

![FIG. 1: The Kerr singular ring and the Kerr congruence formed by oriented twistor null lines and covering the Kerr-Schild spacetime twice.](image)

Twist of the congruence determines the complicate form of the Kerr solution, in spite of the extremely simple form of the metric [1]. One sees also the very important fact that the Kerr-Schild space-time is two-sheeted: the outgoing rays of the Kerr congruence with tangent vector $k^\mu_+ \text{ and the ingoing ones } k^\mu_-$ are positioned on the different sheets of the space and do not interact with each other. The Kerr-Schild formalism [9] takes into account that electromagnetic solutions have to be aligned only to one of two congruences, i.e. the time orientations of the electromagnetic and gravitational fields have to be matched. It meets the requirements of holographic principle [7], since space is projected by the twistor light-like rays on the disk $r = 0$ (“Radial”) coordinate of the Kerr-Schild oblate spheroidal system) which separates the positive and negative sheets of the Kerr space, and also corresponds to the presumable properties of a quantum black hole space-time [8], allowing to separate the in- and out-vacua and to consider the process of evaporation as a scattering of the in-vacuum on black hole.

Elasticity of the horizon [12] follows from the form of function $H$, [9],

$$H = \frac{mr - |\psi|^2/2}{r^2 + a^2 \cos^2 \theta}, \quad (3)$$

where the function $\psi(Y)$ determines electromagnetic field. In the Kerr-Schild class of the exact stationary solutions, function $\phi$ may be any holomorphic function of the complex angular coordinate $Y$. [37]

II. EXACT KERR-SCHILD SOLUTIONS WITH SINGULAR BEAMS

The famous Kerr-Newman solution is the simplest solution of the Kerr-Schild class having $\psi = q = \text{const.}$, where $q$ is the value of charge. However, any holomorphic function $\psi(Y)$ yields also an exact solution of this class [9]. It is known that any holomorphic function on sphere, if it is not a constant, should have at least one singular point. Therefore, all the exact Kerr-Schild solutions, except the Kerr-Newman one, acquire one or more lightlike singular beams which are positioned along the lines of Kerr congruence, while the usual harmonic solutions with smooth angular dependence are absent on the Kerr background. Vector potential of electromagnetic field has the general form

$$\alpha = \alpha_\mu dx^\mu = \frac{1}{2} \Re \left[ \left( \frac{\psi}{r + ia \cos \theta} \right) e^3 + \chi d\bar{Y} \right], \quad (4)$$

where $\chi = \int P^{-2} \psi dY$ and $\bar{Y}$ being kept constant in this integration. The expression $d\bar{Y}$ represents gradient of the complex surfaces $Y = \text{const.}$ obeying the conditions

$$Y_{22} = Y_{44} = 0. \quad (5)$$

These surfaces are totally null, spanned by the tetrad forms $e^1$ and $e^3$. Similarly, $d\bar{Y}$ is spanned by $e^2$ and $e^3$, and therefore, $\alpha^\mu$ is spanned by $e^1$, $e^2$ and $e^3$. So, using the null tetrad orthogonality relations $(e^1)^2 = (e^2)^2 = (e^3)^2 = 0$ and $e^1 e^3 = e^2 e^3 = 0$ and $e^3$, one obtains that vector potential satisfies the alignment condition $\alpha_\mu e^{3\mu} = P \alpha_\mu k^\mu = 0$. The ‘elementary’ beams, formed by a single pole $\psi(Y) = q/(Y - \bar{Y})$ at the point $\bar{Y} \in S^2$, propagate along the line of Kerr congruence in direction $k^\mu$ corresponding to $Y = \bar{Y}$. They play exceptional role, turning in the far zone (see [20]) into uniform string-like singular pp-wave (A.Peres) solutions, [10], which have very important quantum properties, being exact solutions in string theory with vanishing all quantum corrections [21].

In general, holomorphic function may be expanded in the Loron series containing a singular part $\psi(Y) = \sum_{n=-\infty}^0 q_n Y^n$, which can be represented by a series of the above poles, and a regular one $\psi(Y) = \sum_{n=0}^\infty q_n Y^n$. We shall see later that the regular polynomial part plays very important role in the nonstationary solutions. It
was shown, that singular light-like beams deform topologically the horizon, forming the holes connecting the internal and external regions of black hole, see [13]. The ‘elementary’ single-pole solution may trivially be extended to the case of arbitrary number of single poles

\[ \psi(Y) = \sum_{i} \frac{q_i}{Y - Y_i}, \tag{6} \]

in different directions \( Y_i = e^{i\phi} \tan \frac{\theta}{2} \). Elementary excitation \( \psi_i(Y) = q_i/(Y - Y_i) \), describes a singular light-like beam (pp-string) along the null ray of the Kerr congruence in direction \( k_i^a = k^a(Y_i, Y_i) \). The vector potential \( \psi \) is trivially generalized to a sum over beams, where for the single \( i \)-th beam \( \chi_i = q_i (P_f^{-2} \ln(Y - Y_i) + \text{const.} + \mathcal{O}(Y - Y_i)) \), and \( P_f = (1 + Y_i Y_i)/\sqrt{2} \). The corresponding vector field \( \mathcal{E} \) gives rise to electromagnetic field \( f = \frac{1}{2} F_{ab} e^a \wedge e^b = -d\alpha \) which is aligned with the Kerr congruence,

\[ \alpha_{ab} k^a = 0, \quad k^a P^a_P = \lambda k^a, \tag{7} \]

where \( \lambda = \text{Re} \left[ \psi/(r + ia \cos \theta)^2 \right] \). It was shown that these beams have strong back reaction on the metric, via function \( \psi_i(Y) \) entering in the multibeam solution \( (a) \) is a particular case of the exact solutions obtained by Debnay, Kerr and Schild (DKS) in seminal work [3]. It should be emphasized that in the usual perturbative approach the beam solutions are absent, because the important alignment condition are dropped out of the perturbative equations, and a mixing of the in and out fields occurs on the same sheet of metric.

The appearance of light-like beam pulses is a pure classical effect, however, it allowed us in [13] to put some conjectures concerning semiclassical treatment of black hole of interaction with weak electromagnetic field and, in particular, with electromagnetic vacuum. Since the black hole horizon is extra sensitive to electromagnetic excitations of black hole, it may also be sensitive to zero point field (ZPF) which is exhibited classically as Casimir effect. Therefore, as it has been discussed by many authors, vacuum fluctuations have to generate fluctuations of the metric and horizon. Note that such point of view assumes tacitly that there exist some semiclassical prequantum geometry which lies beyond the usual classical gravity and describes the back reaction of the vacuum fluctuations to metric.

The exact Kerr-Schild solutions show that the usual smooth harmonic basic functions are absent on the Kerr background, giving the way to singular beams forming an overfilled system of the coherent non-normalizable states. It rises the questions on the finiteness of the energy of the Kerr-Schild beam-like solutions. Considering the singular beams as a fine-grained vacuum structure, we should take into account that the electromagnetic vacuum energy is divergent, and a transfer to the classical Einstein-Maxwell theory demands at least a regularization of the stress-energy tensor. For example, for calculation of the Casimir energy, de Witt uses the regularization by subtraction \( \mathcal{D}_2 \), \( T_{\mu \nu}^{(reg)} = T_{\mu \nu} - \langle 0 | T_{\mu \nu} | 0 \rangle \), under the condition \( T^{(reg)}_{\mu \nu} \gamma_{\mu \nu} = 0 \). As opposed to a cumulative action of the vacuum in Casimir effect, the treatment of the fine-grained structure of fluctuations may require another form of the regularization. As we shall see, the structure of the Debney-Kerr-Schild equations (DKS), [3], suggests a regularization which retains the fluctuating fine-grained structure of the horizon and the semiclassical pre-geometry. To transfer to the usual classical gravity with a smooth background geometry and get the usual scenarios of collapse, an extra regularization of the fine-grained structure is necessary.

The real beams has to be finished at some distant black holes or particles, or at the infinitely distant matter. It is confirmed by the treatment of multicenter Kerr-Schild solutions, [13], in which the similar beams appear between the distant sources, being supported by a twistor line common for them.

The exact pp-wave Peres solutions may have a carrier frequency and a finite extension. Therefore, it is desirable to consider a minimal generalization of the exact stationary Kerr-Schild beam-like solution to time-dependent beam pulses, which is necessary to consider a time-dependent process of scattering.

### III. TIME-DEPENDENT KERR-SCHILD SOLUTIONS

The closest time-dependent generalization to \( (d) \) is given by the form

\[ \psi(Y, \tau) = \sum_{i} c_i(\tau)(Y - Y_i)^{-1}, \tag{8} \]

where one assumes that an elementary beam has \( c_i(\tau) = q_i(\tau)e^{i\omega_i \tau} \) where \( q_i(\tau) \) is amplitude and \( \omega_i \) is carrier frequency. Existence of the exact nonstationary solutions of this type is a priory problematic, and the problem may have two stages: 1) exact solutions for electromagnetic
field on the Kerr background, ii) self-consistent solutions, taking into account the back reaction of the electromagnetic field on metric. General equations for nonstationary electromagnetic Kerr-Schild solutions were obtained by Debney, Kerr and Schild (DKS) in 1968, and the general electromagnetic solution on the Kerr-Schild background was obtained only in 2004. The analyzed in wave solutions contained singular beams along the ± axes, and there appeared the conjecture that, for certain exclusions, the exact nonstationary electromagnetic solutions on the Kerr background should acquire the beams. Similar to stationary case, the nonstationary Kerr-Schild solutions contain the singular beams which change the structure of black hole horizon. The considered below important peculiarities of the nonstationary solutions are determined by a specific structure of the DKS equations. Electromagnetic field is described by the self-dual tetrad components,

\[ \mathcal{F}_{ab} = AZ^2, \quad \mathcal{F}_{31} = \gamma Z - (AZ)_1, \]

where \( \mathcal{F}_{ab} = e^a_k e^b_l \mathcal{F}_{kl} \), and the function \( Z \) is a complex expansion of the congruence, \( Z = Y_1 \). For the Kerr-Newman solution at rest \( Z \) is inversely proportional to a complex radial distance \( Z = -P/(r + ia \cos \theta) \). Here \( P \) is a conformal factor which is determined by Killing vector of the solution. For a black hole at rest \( P = 2^{-1/2} (1 + YY) \). The equations for function \( A \) are the same as for stationary case \( (AP^2)_{,\alpha} = 0, A_{,4} = 0 \). They have the general solution \( A = \psi P^{-2} \), where function \( \psi \) has to satisfy

\[ \psi_{,2} = \psi_{,4} = 0. \]

So, in the nonstationary case, there appears the unique difference that function \( \psi \) has to depend on extra retarded-time parameter \( \tau \), which must obey the equations \( \tau_{,2} = \tau_{,4} = 0 \). The principal difference from the stationary case is contained in the second electromagnetic DKS equation which in was reduced to very simple form

\[ (\gamma P)_{,\psi} = -\dot{A}, \]

showing that any nonstationarity in electromagnetic field \( \dot{A} \neq 0 \) generates an extra function \( \gamma \) which, in accord with [9], generates also the lightlike electromagnetic radiation along the Kerr congruence. Such a radiation is well-known for the Vaidya ‘shining star’ solution, in which the field \( A = \psi P^{-2} \) is absent and \( \gamma \) is incoherent, being related with the loss of total mass into radiation,

\[ \dot{m} = -\frac{1}{2} P^2 \gamma \gamma. \]

This is one of two gravitational equations determining self-consistency of the Kerr-Schild solution. The most important consequence following from DKS equations in the nonstationary case is the fact that the field \( \gamma \) appears inevitable, however it does not contribute to deformation of the horizon, since it is absent in the function \( H \) of [3]. Its back reaction on metric is smooth and circumstantial, acting only via the slowly decreasing mass parameter \( m \).

Therefore, in the nonstationary Kerr-Schild case we obtain that the lightlike fields, determined by functions \( \psi \) and \( \gamma \), have essentially different impact on the horizon.

Function \( \psi = \psi(\tau, Y) \) obeys the equation [10] which shows that the retarded time \( \tau \) has to satisfy the conditions similar to [3], and therefore, gradient of \( \tau \) has to be aligned to congruence,

\[ k^\mu \tau_{,\mu} = 0. \]

It was obtained in [20] that the corresponding retarded-time parameter has the form

\[ \tau = t - r - ia \cos \theta, \]

and the general solution of the equations [11] takes the form

\[ \gamma = \frac{2^{1/2} \psi}{P^2 Y} + \phi(Y, \tau)/P, \]

where \( \phi(Y, \tau) \) is the second arbitrary analytic function of \( Y \) and \( \tau \) which is solution of homogenous equation [11] corresponding to \( \dot{A} = \psi = 0 \).

Since function \( \phi \) contributes only to \( \gamma \) it does not impact on the form of the horizon too. Important role of this function is obtained from analysis of the second gravitational Kerr-Schild equation [eq.(5.44) in [9]] which is reduced to the form (see [13])

\[ m_{,\psi} = \psi \gamma P. \]

If we note that \( \gamma \sim \psi \approx \omega \psi \), we obtain that the r.h.s. of this equation tends to zero in the low-frequency limit, as well as the r.h.s of the equation for \( \dot{m} \). So, the full solutions tend to exact ones (consistent with gravity) in the low-frequency limit [13, 14].

However, there is a more consequent way for interpretation of these solutions which is close to a quantum version of the Einstein-Maxwell equations. So far as we consider the vacuum electromagnetic field, one has to assume that electromagnetic field should have an operator meaning and a regularization of the stress-energy (r.h.s. of the gravitational DKS equations) is necessary.

Note that the Vaidya ‘shining star’ solution [10] corresponds to \( \psi = 0 \), and the equation [16] shows that \( m \) is independent from \( Y \), while the field \( \gamma \) is assumed to be incoherent and has to be considered with averaged r.h.s., \( \frac{1}{2} \frac{< P^2 \gamma \gamma >}{2} \). This approach can be extended to the r.h.s of the both gravitational equations which acquire the form

\[ m_{,\psi} = \frac{< P \psi \gamma >}{2}, \]

and may be considered as a semiclassical analog of quantum approach. [39]
Since function $\phi$ is free, its form and parameters (positions of poles) of the function $\phi/P$ may be tuned to cancel the poles of function $\hat{\psi} = \sum_i \hat{c}_i(\tau)/(Y - Y_i)$ in $\mathbb{R}^4$. We set

$$\phi_i^{(\text{tun})}(Y, \tau) = -\frac{2^{1/2} \hat{c}_i(\tau)}{Y(Y - Y_i) P_i},$$

where

$$P_i = P(Y, \bar{Y}_i) = 2^{-1/2}(1 + Y \bar{Y}_i)$$

is analytic in $Y$, which provides required analyticity of $\phi_i^{(\text{tun})}(Y, \tau)$. Using the equality

$$(P_i - P)/(Y_i - Y) = \frac{Y (Y_i - \bar{Y})}{\sqrt{2} (Y_i - Y)},$$

we obtain that the regularized field

$$\gamma(\text{reg}) = \frac{2^{1/2} \hat{\psi}}{P^2 Y} + \sum_i \phi_i^{(\text{tun})}(Y, \tau)/P$$

takes the form

$$\gamma(\text{reg}) = \frac{1}{P^2} \sum_i \hat{c}_i \bar{Y}_i - \bar{Y}.$$ (22)

It should be mentioned that $\gamma(\text{reg})$ is exact regular solution of (11). Therefore, the exact solutions of DKS equation (11) contains remarkable subclass of self-regularized solutions $\gamma(\text{reg})$.

R.h.s. of the equation (16) takes the form

$$\psi\tilde{\gamma}(\text{reg})P = \sum_{i, k} \psi_i \tilde{\gamma}_k(\text{reg})P = \sum_{i, k} c_i \hat{c}_k (Y - Y_k)/PP_k(Y - Y_i)/(Y - Y_k),$$

where the terms with $i \neq k$ are not correlated and drop out after averaging. The rest is

$$< \psi\tilde{\gamma}(\text{reg})P > = - \sum_k \frac{c_k \hat{c}_k}{PP_k(Y - Y_k)}.$$ (23)

Equation (16) may be integrated using the Cauchy integral formula, and we obtain the expression

$$< m >_t = m_0 - 2\pi i \sum_k \frac{c_k \hat{c}_k}{|P_k|^2}$$

containing contributions from the residues at singular points $Y$.

When $c_i(\tau)$ is expressed via amplitudes $q_i(\tau)$ and carrier frequencies $\omega_i$, $c_i(\tau) = q_i(\tau)e^{-i\omega_i \tau}$, the impact of the carrier frequencies disappears,

$$< m >_t = m_0 + 2\pi \sum_k \omega_k \sum_k < \frac{q_k \hat{q}_k}{|P_{kk}|^2} >,$$ (25)

however, the mass term retains slow fluctuations and an angular non-homogeneity caused by casual distribution of the beams in different angular directions.

It is known that the second gravitational equation in (17) is really a definition of the loss of mass in radiation. The time averaging removes again the terms with $i \neq k$ and yields

$$< \dot{m} >_t = - \frac{1}{2} \sum_k \frac{\hat{c}_k \dot{q}_k}{P^2 |P_k|^2}.$$ (26)

In terms of the amplitudes of beams we obtain

$$< \dot{m} >_t = - \frac{1}{2} \sum_k \omega_k^2 < \frac{\hat{q}_k \dot{q}_k}{|P_{kk}|^2} >,$$ (27)

which shows influence of the single beam to the mass evaporation. These sporadic fluctuations of the mass term caused by the individual beam pulses may also be averaged over the time and angular directions, however, this smoothing does not remove the sharp back-reaction of the beams to metric and horizon produced by poles in $\psi(Y, \tau)$ in agreement with (3).

Therefore, the obtained solutions are exact for the time-dependent electromagnetic field on the Kerr-Schild background and, up to our approximation which neglects the recoil, the obtained solutions are consistent with the Einstein-Maxwell system of equations with averaged stress-energy tensor.

**IV. HOLOGRAPHIC GRAVITY OF THE TIME-DEPENDENT KS SOLUTIONS**

The obtained very broad class of semi-classical solutions to Einstein-Maxwell equations have two principal peculiarities: the solutions are time-dependent and contain a sporadic flow of beam pulses. Therefore, the solutions reveals a classical fine-grained fluctuating geometry of twistor-beams and exhibit a mechanism of evaporation provided by the specific structure of DKS equations. They show explicitly that the outgoing radiation contains two components determined by functions $\psi$ and $\gamma = \gamma(\text{reg})$. Both the functions are creating the null electromagnetic radiation along the twistor null lines of the Kerr congruence in agreement with (1), however, they play their own specific role:

a) function $\psi$ describes a casual set of the outgoing singular beams which perforate the horizon, forming fluctuating micro-holes breaking impenetrability of the horizon.

b) the field determined by function $\gamma(\text{reg})$ is regular and akin to the Vaidya ‘shining star’ radiation. In the agreement with $\dot{m} = -\frac{1}{2} P^2 \gamma(\text{reg})\tilde{\gamma}(\text{reg})$, it is responsible for the mass evaporation.

Therefore, the resulting evaporation represents a classical analog of the quantum tunnelling process.
Holographic interpretation of the mysterious twosheetedness of the Kerr-Schild geometry allows one to treat evaporation as a scattering and reveals the important role of the ‘negative’ sheet of the Kerr geometry, showing that Kerr’s twosheetedness represents a classical analog of the required holographic structure of a quantum black-hole space-time predicted by Stephens, t’ Hooft and Whiting (StHW) in 
18. The holographic StHW space-time is to be divided into two causally-related ‘in’ and ‘out’ regions joined by a 2+1 dimensional (shell-like) boundary which is holographically dual to the ‘in’ and ‘out’ bulk regions. Twosheeted structure of the Kerr geometry is perfectly adapted to holographic StHW structure. It has to be unfolded, forming the in and out bulk regions separated by the Kerr BH source, as it is shown on the Penrose conformal diagram, Fig.3.

![Penrose Conformal Diagrams](image)

**FIG. 3**: Penrose conformal diagrams. Unfolding of the auxiliary $M^4$ space of the Kerr spacetime to StHW structure of a quantum BH spacetime.

The BH appears as a holographic image created by the projection performed by twistor rays from past infinity $I^-$. After scattering on the Kerr BH, the out-going beams propagate along the twistor rays of Kerr congruence to future infinity $I^+$. Therefore, radiation from BH appears as a result of scattering of the in-vacuum field on the shell-like source of the Kerr BH which may be, in particular, a shell of collapsed matter positioned very close to horizon, or a matter already collapsed under the horizon of an even BH, either also a disk-like source of the Kerr-Newman geometry without horizon [23, 24, 25, 26, 27], as it was considered in the models of the Kerr spinning particle. Similar, authors of [23] argue that such a source may also be positioned inside of the horizon and the existence of the horizon, in principle, is not important for derivation of the BH radiation. In all the cases the Kerr source may be replaced by a rotating shell which separates the ‘in’ and ‘out’ bulk regions of the Kerr space-time. In the simplest case it is the disk $r = 0$ spanned by the Kerr singular ring - the Israel-Hamity Kerr’s source [23]. Since the twistor rays of the Kerr congruence are time oriented null lines, the Kerr-Schild alignment condition (7) plays specific role of time ordering on the considered semi-classical level. However, the Kerr-Schild solutions are analytic solutions which may be extended analytically back onto negative sheet of the Kerr space, and the treatment of the StHW Kerr-Schild structure together with the holographic conception of scattering, assume that this analyticity is broken by the shell-like Kerr source. It allows us to propose that the in-sheet, $r < 0$, is to be residence of the vacuum in-fields $\gamma_{in}$. In some sense this sheet may be considered as nonphysical one, since the black hole horizon is absent by $r < 0$. One can also suppose that the in-field $A_{in}$ is also absent there (at least, we have to put $A_{in} = 0$), which yields $\gamma_{in} = -\gamma_{(tun)}$. Thus, we obtain that the ‘negative’ in-sheet is a region of the homogenous solutions of DKS equations which may contain singular beams but do not produce singular back reaction to the metric. Therefore, the metric turns out to be smooth by $r < 0$. The scattering of the field $\gamma_{in}$ at the Kerr source creates the outgoing electromagnetic singular beams $A_{out}$, accompanied by singular deformations of the metric and horizons and by the regular outgoing incoherent thermal field $\gamma_{out} = \gamma_{reg}$.

The holographic bounce conception [8], related with the treatment of the BH evaporation as a scattering, has the specific feature that the beams are scattered on the 2+1 dimensional shell-like source. Taking into account, that the solutions on the in- and out-Kerr’s sheets have in this case their own doubling by analytic extensions to another half-sheet, we arrive at a possible interplay of the four half-sheets. Moreover, from the treatment of multi-center Kerr-Schild solutions [17], we know that any extra matter source adds two more sheets to the KS geometry. Therefore, the presence of the matter shell source adds at least two more sheets to the KS space-time and have to lead to the KS solutions with a shell-like singularity. Therefore, the usual Greene functions have to be generalized to some KS solutions which are singular at the shell. It is clear that topological interplay of these half-sheets may be related with the origin of nontrivial commutation relations. These questions need especial treatment which goes out of the frame of this work.

The next important problem is related with spectrum of the evaporation. The obtained analytical solutions have continuous spectrum, displaying that the discrete quantum spectrum can appear only as a result of interaction of the in-vacua with a matter source of the BH.

The simplest source of the even Kerr BH is the disk positioned at $r = 0$. In fact, it has a stringy structure, representing a light-like closed Kerr string positioned on the boundary of the Kerr disk [18, 23, 26, 28, 29]. This string is reminiscent of the fundamental heterotic string [30], the Sen solution to low energy string theory [31].

The averaged stress-energy tensor of the Kerr-Schild solutions,

$$T^\mu_\nu = Z \bar{Z} \langle \gamma_{(reg)} \bar{\gamma}_{(reg)} \rangle <k^\mu k^\nu>, \quad (28)$$

is determined by term $\gamma_{out} = \gamma_{(reg)}$ and is independent from the position and even from the presence of the hori-
zon. It has axial symmetry and the symmetry of time-translations which characterize the symmetry of a conformal field theory determined by the function $\gamma_{out}(Y, \tau)$. For the Kerr geometry at rest $Z = P/(r + ia \cos \theta)$. Considering the disk $r = 0$ as a holographically dual boundary, we obtain that the stress-energy tensor turns out to be singular in the equatorial plane $\theta = \pi/2$, which corresponds to the CFT on the Kerr singular ring, $r = \cos \theta = 0$, i.e., CFT of the Kerr light-like string \cite{29, 30}. Therefore, the quantum modes of the excitations (in this case only one side movers) appear as a result of resonance of the in-vacuum field on the singular string-like source of the Kerr BH. It gives evidences that the obtained radiation of Vaidya type corresponds to Hawking radiation in the agreement with two-dimensional CFT of the string theory related with the Kerr singular ring. One can chose $\partial_\phi$ and $\partial_t$ as generators of a Virasoro algebra for the spacetime diffeomorphisms associated with CFT on the cylinder $(\phi, t)$, and following \cite{3, 4} use the Cardy formula for spectrum. \cite{3}. Formally, it is similar to the given by Strominger derivations for the boundary CFT positioned at infinity, but we have absolutely different physical mechanism related with excitations of the Kerr string.

In the many recent works, for example \cite{5, 6, 7}, the holographically dual shell source represents a dense matter concentrated near the horizon of the Kerr BH. All the above features of the bounce StHW model are retained for exclusion of the clear exhibition of the stringy structure of the Kerr source. The horizon of rotating BH takes the form of a rigidly rotating disk, oblateness of which depends on the velocity of angular rotation, and a string-like region is formed on the edge rim of the disk only for the quickly rotating extremal black-holes. It all other cases the above sharp stringy effect turns out to be smeared, and presumably, spectrum may be formed as a mixing of spectrum for many strings.

Details of the quantum version of the holographic Kerr-Schild space and the thermal processes related with the quantum resonance demand extra analysis and have to be considered elsewhere.

It should also be noted that in the obtained solutions we neglected by the recoil, presence of which essentially complicates DKS equations and leads also to extra problems with geometrical matching of the in and out sheets. There appears also a lot of other important and interesting questions which we have to leave for further considerations.

V. CONCLUSION

The obtained Kerr-Schild solutions and the holographic representation of the Kerr twosheetedness showed that the Kerr-Schild geometry is a semiclassical analog of the holographic StHW space-time suggested for quantum BH background. The obtained solutions are time-dependent and describe a fine-grained fluctuating holographic BH geometry which takes an intermediate position between the classical and quantum gravity. Basic elements of this geometry are beam pulses which take asymptotically the form of pp-waves. Singular lines of the beams are supported by twistor null lines of the Kerr-Schild geometry. This fine-grained structure of spacetime formed by twistor beams represents a holographic alternative to the considered in lattice structure of vacuum formed by giant masses positioned at the nodes of the lattice.

Therefore, along with the old Penrose statements on the principal role of twistors in the structure of space-time, there appear extra evidences that the oriented twistors may be considered as basic elements of the vacuum structure. Principal advantage of the twistor vacuum structure with respect to the usual lattice structure is that the vacuum based on twistors possesses the explicit Lorentz invariance and explicit time orientation. Twistor description of the massless fields based on the WZW or topological B-model represents the basis of the Nair-Witten concept on the scattering of the gauge amplitudes in twistor space \cite{34, 35}. The Penrose claim that twistors may be considered as primary objects of 4D space-time with respect to the space-time points (which may be considered as geometrically dual to twistors) is explicitly illustrated by the holographic structure of the Kerr-Schild black hole space-time which allows one to describe the structure of the Kerr black-hole via the Kerr congruence of twistor null lines.

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[38] We neglect the recoil, \(\dot{P} = 0\), assuming that that the mass of black hole is much greater then the energy of excitation.
[39] In [32] the general Kerr-Schild solution (15) is applied to a black hole-disk system and some interesting properties of the averaged stress-energy tensor are observed.

\[Y(x) \text{ determines also the Kerr-Schild tetrad } e^a: \]
\[e^1 = d\zeta - Y dv, \quad e^2 = d\bar{\zeta} - \bar{Y} dv, \quad e^3 = P k_\mu dx^\mu, \quad e^4 = dv + he^3, \quad h = H P^{-2}, \quad \bar{P} = (1 + Y \bar{Y})/\sqrt{2}.\]