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On Phase Transitions during Collisions Near the Horizon of Black Holes

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Abstract: During particle collisions in the vicinity of the horizon of black holes, it is possible to achieve energies and temperatures corresponding to phase transitions in particle physics. It is shown that the sizes of the regions of the new phase are of the order of the Compton length for the corresponding mass scale. The lifetime is also on the order of the Compton time. It is shown that the inverse influence of the energy density in the electro-weak phase transition in collisions on the space–time metric can be neglected.

Keywords: black hole; symmetry breaking; phase transitions

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1. Introduction

The works of A.A. Friedman [1,2], written 100 years ago, in which solutions were obtained for an expanding homogeneous isotropic universe [3], are the theoretical basis of the modern standard cosmological model. The discovery in 1965 of relic radiation [4,5] indicates that in the model of the expanding early Universe, there were times when the temperature of matter was so high that phase transitions predicted by the theory of elementary particles could occur. There are three such phase transitions in the standard model of particle physics [6–8]:

(1) Between quark–gluon plasma and hadrons at the energies E of the order of 200 MeV. The corresponding temperature $T = E/k_B \approx 10^{12}$ K, where $k_B \approx 1.38 \cdot 10^{-23}$ J/K is the Boltzmann constant, may have taken place in the expanding Universe during the order of 10^{-6} s after the Big Bang.
(2) An electro-weak phase transition at energies of the order of $E_W \approx 100$ GeV. The corresponding temperature $T_W \approx 10^{15}$ K could have taken place during the order of 10^{-12} s after the Big Bang.
(3) The grand unification phase transition at energies.

$E_{\rm GUT} \approx 10^{16}$ GeV. The temperature corresponding to the energy of the grand unification phase transition $T_{\rm GUT} = E_{\rm GUT}/k_B \approx 10^{29}$ K may not have been achieved in the early universe in models with an inflationary stage in which the heating temperature is significantly lower than $T_{\rm GUT}$. In models with a radiation-dominant stage in the early Universe, the temperature $T_{\rm GUT}$ could be reached at times of the order of 10^{-38} s.

The study of the properties of matter at such temperatures and the phenomena at these phase transitions is of undoubted theoretical interest. Is it possible to achieve such temperatures in experiments on the Earth? The maximum high temperature for macroscopic parts of the substance is achieved at the time of nuclear explosion and can be on the order of 10^8 K. This is significantly less than the temperature of even the quark–gluon phase transition [9].

The highest temperature achieved in experiments on Earth refers to microscopic quantities of matter and is obtained when heavy element nuclei collide in particle accelerators. A temperature of $4 \cdot 10^{12}$ K was obtained from a collision of gold nuclei in Brookhaven National Laboratory (United States) in 2010 [10]. In 2012, it was reported that a temperature of $5 \cdot 10^{12}$ K was reached when the lead nuclei collided at the Large Hadron Collider [11]. At such temperatures, hadron matter transforms into the quark–gluon plasma state. However, such temperatures are more than two orders of magnitude less than the temperature of the electro-weak phase transition.

Thus, macroscopic amounts of matter in the state of phase transition of elementary particle physics in laboratories on Earth cannot be obtained, and microscopic amounts can be obtained only for the phase transition in the quark–gluon plasma state.

Is it possible to observe matter at the temperatures of the phase transitions of particle physics in astrophysical processes at present? Brightly luminous accretion discs formed when matter falls into black holes have a visible temperature of hundreds of millions of Kelvin degrees [12]. As shown in our work [13], in the processes of collisions of particles near the horizon of black holes, it is possible to achieve energies in the system of the center of mass of colliding particles on the order of the energy scale of the electro-weak phase transition? A summary of these results is presented in Section 2.

Here we will consider questions about the size of the regions of the phase transition region obtained in a collision and the lifetime of such a region. To do this, in Section 3, we apply formulas for the energy density and radiation intensity of a gas of relativistic particles. The possibility of obtaining an electro-weak phase transition in a macroscopic volume during a collision in the vicinity of supermassive black holes is studied in Section 4. The influence of the matter energy–momentum tensor in the phase transition region on the space–time metric will be evaluated in Section 5.

2. The High-Energy Collisions Near the Horizon of Black Holes

The Kerr metric of a rotating black hole [14] in the Boyer–Lindquist coordinates [15] has the following form:

$$ds^2 = \frac{\rho^2 \Delta}{\Sigma^2} c^2 dt^2 - \frac{\sin^2\theta}{\rho^2} \Sigma^2 (d\varphi - \omega dt)^2 - \frac{\rho^2}{\Delta} dr^2 - \rho^2 d\theta^2, \tag{1}$$

where

$$\rho^2 = r^2 + \frac{a^2}{c^2}\cos^2\theta, \quad \Delta = r^2 - \frac{2GMr}{c^2} + \frac{a^2}{c^2}, \tag{2}$$

$$\Sigma^2 = \left(r^2 + \frac{a^2}{c^2}\right)^2 - \frac{a^2}{c^2}\sin^2\theta\, \Delta, \quad \omega = \frac{2GMra}{\Sigma^2 c^2}, \tag{3}$$

G is the gravitational constant, c is the speed of light and M and aM are the mass and angular momentum of the black hole, respectively. We accept that $0 \leq a \leq GM/c$. The event horizon of the Kerr black hole has the radial coordinate

$$r = r_H \equiv \frac{G}{c^2}\left(M + \sqrt{M^2 - \left(\frac{ac}{G}\right)^2}\right). \tag{4}$$

According to [16], the squared energy of a collision of two particles with a mass m with the angular momenta L_1 and L_2 in the center-of-mass system, which are nonrelativistic at infinity and are freely incident on a black hole with the angular momentum aM, is given by the expression

$$\frac{E_{\text{c.m}}^2}{m^2c^4} = \frac{2}{x(x^2-2x+A^2)}\Bigg[2A^2(1+x) - 2A(l_1+l_2) - l_1 l_2(x-2) + 2(x-1)x^2 \\ - \sqrt{2(A-l_2)^2 - l_2^2 x + 2x^2}\sqrt{2(A-l_1)^2 - l_1^2 x + 2x^2}\Bigg], \tag{5}$$

where $x = rc^2/GM$, $l = Lc/GmM$ and $A = ac/GM$. Expression (5) has a singularity on the event horizon. In the general case, the limit value of the collision energy for two particles with masses m_1, m_2, energies E_1, E_2 and angular momenta J_1, J_2 is

$$E_{\text{c.m}}^2(r \to r_H) = \frac{c^6(J_{1H}J_2 - J_{2H}J_1)^2}{G^2M^2(J_{1H}-J_1)(J_{2H}-J_2)} + m_1^2c^4\left[1 + \frac{J_{2H}-J_2}{J_{1H}-J_1}\right] + m_2^2c^4\left[1 + \frac{J_{1H}-J_1}{J_{2H}-J_2}\right], \tag{6}$$

where $J_{nH} = 2E_n r_H/A$. If the angular momentum of one of the particles tends to J_{nH}, then the expression for the energy (6) diverges. This is the so-called Banados–Silk–West effect. Note that despite the unlimited increase in collision energy in the center of mass system, the energy that can be extracted at a large distance from a black hole cannot exceed $E_1 + E_2$ (assuming no Penrose effect [17]). This follows from the law of energy conservation.

A particle having a critical angular momentum value can travel from infinity to the event horizon of a black hole only in the case of an extremely rotating black hole $A = 1$. In other cases, particles with large angular momentum values are prevented from falling onto the horizon by the potential barrier of the effective potential. As shown in [18,19], the super high center-of-mass energy can be achieved in multiple collisions near nonextreme black holes. To reach the horizon, particles incident from infinity should have an angular momentum low in absolute value. The angular momentum of one of the particles necessary for a high-energy collision can be acquired either in multiple collisions or in the interaction with the electromagnetic field of the accretion disk. A similar effect for electrically charged black holes was discovered in [20]. Real astrophysical black holes are surrounded by matter (for example, they have an accretion disk). The possibility of particles colliding with unlimited energy near the horizon of such "dirty" black holes also takes place [21].

The value of the collision parameters corresponding to the temperature of the elementary particles phase transitions may depend on the type of black holes. In the case of Kerr black holes, the estimates for the distance from the horizon, where the collision energies required for phase transitions of elementary particles, can be achieved are given in our work [13]. So, for elementary particles with a mass m, the value of the temperature T is reached near the extreme rotating black hole at the distance

$$r - r_H \approx 2r_H\left(\frac{mc^2}{k_BT}\right)^2. \tag{7}$$

For the proton mass, the electro-weak temperature can be reached at the distance $r - r_H = 2\times10^{-4}r_H$. This amounts to tens of centimeters for stellar-mass black holes and hundreds of thousands of kilometers for supermassive black holes. In the mechanism of multiple collisions near the horizon of (not extreme) rotating black holes, such temperatures can be achieved at larger distances [13]. Therefore, collisions in which phase transition temperatures are reached can, in principle, occur in the vicinity of stellar-mass black holes for elementary particles and in the case of supermassive black holes for macroscopic bodies.

Next, we estimate the size of the phase transition region and the lifetime of the state with the new phase.

3. Size and Lifetime of the New Phase

The energy density of the photon gas can be calculated by the known formula: (see (63,14) in [22])

$$\varepsilon = \frac{4\sigma}{c}T^4, \tag{8}$$

where σ is the Stefan–Boltzmann constant,

$$\sigma = \frac{\pi^2 k_B^4}{60\hbar^3 c^2} \approx 5.67 \cdot 10^{-8} \frac{\mathrm{W}}{\mathrm{m}^2 \cdot \mathrm{K}^4}, \tag{9}$$

and $\hbar$ is the reduced Planck constant. At ultra-high temperatures, other elementary particles should also contribute to the energy density of matter. Their contribution is taken into account using the factor g_{eff}, which describes the number of effective massless degrees of freedom of particles of the standard model of particle physics.

$$\varepsilon = g_{\text{eff}} \frac{\pi^2 k_B^4}{30\hbar^3 c^3} T^4 = g_{\text{eff}} \frac{2\sigma}{c} T^4. \tag{10}$$

Under this definition, the photon's contribution to g_{eff} is two, according to the photon's two polarization states. In the general case, one has the following [6]

$$g_{\text{eff}} = \sum_{i=\text{bosons}} g_i \left(\frac{T_i}{T}\right)^4 + \frac{7}{8} \sum_{i=\text{fermions}} g_i \left(\frac{T_i}{T}\right)^4. \tag{11}$$

Here, it is assumed that the equilibrium temperature T_i of particles of type i may differ from T. For example, in the Universe at present the temperature of cosmic microwave background radiation is equal to 2.7 K, and estimates for the temperature of the neutrino gas give 1.95 K. Photon gas after the moment of the last collisions of the cosmological neutrinos with cosmological plasma at energies of 2–3 MeV was still heated up in the annihilation process of cosmological positrons with electrons.

The value of g_{eff} in the standard model of particle physics depends on temperature. For T in the interval 1 MeV $< T <$ 100 MeV, which takes neutrinos into account, electrons and positrons lead to $g_{\text{eff}} = 10.75$. At temperatures above 300 GeV, all standard model particles (photons, $W^\pm$, Z^0 bosons, eight gluons, three generations of quarks and leptons and the Higgs boson) must contribute to (10), which leads [6] to the value of $g_{\text{eff}} = 106.75$. The graph of g_{eff}, which depends on temperature, is presented in [6] on page 65, Figure 3.5.

Denoting $k_B T = mc^2$, where m is the characteristic mass scale, we obtain from (10) for the energy density of radiation of all types of particles

$$\varepsilon = g_{\text{eff}} \frac{\pi^2 m^4 c^5}{30\hbar^3} = g_{\text{eff}} \frac{\pi^2}{30} \frac{mc^2}{l_C^3}, \tag{12}$$

where $l_C = \hbar/mc$ is the (reduced) Compton wavelength corresponding to the mass m. The pressure corresponds to a value three times less

$$p = \frac{\varepsilon}{3} = g_{\text{eff}} \frac{\pi^2}{90} \frac{mc^2}{l_C^3}. \tag{13}$$

The size R_0 of the area in which the heated drop of a new phase of matter can form after a collision is estimated from the relation

$$E_{\text{c.m.}} = \frac{4}{3}\pi R_0^3 \varepsilon, \tag{14}$$

It is assumed that the region of the new phase is a sphere with the radius R_0. Then, one obtains

$$R_0 = \frac{l_C}{\pi} \sqrt[3]{\frac{45}{2g_{\text{eff}}} \frac{E_{\text{c.m.}}}{mc^2}}. \tag{15}$$

Assuming that the collision energy is of the order of magnitude $E_{\text{c.m.}} \sim g_{\text{eff}}\, \mathrm{mc}^2$, we find that the size of the region phase transition is of the order of the Compton wavelength l_C for a particle of mass m.

Let us estimate the lifetime of a drop of a new phase formed as a result of a collision, generalizing the formula for the radiation intensity of the black body to the case of the presence of additional degrees of freedom described by the quantity g_{eff}

$$J = g_{\text{eff}} \frac{\pi^2 k_B^4}{120\hbar^3 c^2} T^4 = \frac{g_{\text{eff}}}{2} \sigma T^4. \tag{16}$$

Let us write the energy balance equation for an infinitesimal time interval dt

$$d(\varepsilon V) = -JSdt, \tag{17}$$

where V is the volume of new phase drop and S is its surface area. When obtaining estimates by the order of magnitude, we assume that the drop is spherical, and the radius may depend on time due to expansion into the surrounding space. We also assume that during the life of a drop of a new phase, thermodynamic equilibrium takes place in it, and, therefore, we can talk about the temperature of the entire drop, the dependence of temperature on time and use formulas for the equilibrium state of the corresponding relativistic gas. Then, from (17), we obtain

$$\frac{R}{3} d\varepsilon = -(Jdt + \varepsilon dR). \tag{18}$$

Using (10) and (16), we obtain the equation

$$\frac{16}{3} \frac{R}{c} \frac{dT}{T} = -\left(1 + \frac{4}{c}\frac{dR}{dt}\right) dt. \tag{19}$$

By integrating this equation, we obtain

$$T(t) = T(t_0) \exp\left[-\frac{3c}{16} \int_{t_0}^{t} \left(1 + \frac{4}{c}\frac{dR}{dt}\right) \frac{dt}{R}\right]. \tag{20}$$

If the drop radius does not change, i.e., $R \approx R_0 = \text{const}$, then the solution is

$$T(t) = T(t_0) \exp\left[-\frac{3}{16} \frac{c(t - t_0)}{R_0}\right]. \tag{21}$$

Thus, the temperature decreases exponentially, and the lifetime of the new phase is of the order $\tau \approx R_0/c$. Since, according to the Equation (15), the size of the new phase region is assumed to be Compton, the lifetime corresponds to Compton time $\tau_C = \hbar/(\mathrm{mc}^2)$ for a particle of a mass m, corresponding to the phase transition energy.

Taking into account the possible expansion of the area of the new phase, the lifetime of the new phase can only decrease. Let us give formulas under the assumption of a constant expansion rate $dR/dt = v \approx \text{const}$. Then,

$$R(t) = R_0 + v(t - t_0) \tag{22}$$

and, after the integration of (19), we obtain a dependence of the temperature of the region with a new phase on the time as follows

$$T(t) = T(t_0) \left(1 + \frac{v(t - t_0)}{R_0}\right)^{-\frac{3}{16}\left(4 + \frac{c}{v}\right)}. \tag{23}$$

In the limit $v/c \to 0$, one obtains the expression (21).

Thus, the lifetime of the new phase obtained in a collision of elementary particles has the order of Compton time $\hbar/(mc^2)$ for a particle of the characteristic mass scale m. For the quark–gluon phase transition, this time is $\tau \approx 3 \cdot 10^{-24}$ s. For the electro-weak phase transition, this time is $\tau \approx 7 \cdot 10^{-27}$ s.

4. Phase Transition in Macroscopic Volume

To perform a quark–gluon or electro-weak phase transition in a macroscopic volume, it is necessary to collide with ultra-relativistic energies of macroscopic amounts of matter. When ordinary macroscopic bodies collide with such energies, the regions of the new phase can make up a macroscopic volume only if the density of the bodies is comparable to the Compton density characteristic of the phase transition of the mass m (see (12)). Only in this case, the lifetime of the new phase can significantly exceed the Compton time τ_C. Such density of matter occurs only in neutron stars. Collisions of macroscopic objects with ultra-relativistic velocities are possible in the vicinity of the horizon of extremal rotating black holes [13]. The collision of compact objects with star masses near supermassive black holes was considered in [23].

When falling towards the event horizon of a black hole, macroscopic bodies can be destroyed by tidal gravitational forces. Let us estimate the mass of black holes in which it is possible to fall to the event horizon of neutron stars without destruction by tidal forces. For evaluation, we assume that a star is destroyed if the tidal forces for the points of the center of mass and the surface exceed the force of attraction of the points of the surface to the center of the falling body. Let us assume that the falling object (neutron star) is a uniform ball of a density ρ and radius R. Also, let us consider only the nonrotating black hole and radial tidal forces. Then the condition for falling to the horizon without destruction has the form

$$\frac{2GM}{r_g^3} R < \frac{G 4\pi\rho R^3}{3R^2} \tag{24}$$

or (after simple transformations)

$$M > \frac{c^3}{4G^{3/2}}\sqrt{\frac{3}{\pi\rho}}, \quad \frac{M}{M_\odot} > 1.9 \cdot 10^8 \sqrt{\frac{\rho_w}{\rho}}, \tag{25}$$

where $M_\odot$ is the Sun mass, $\rho_w = 10^3\,\mathrm{kg/m^3}$ is the water density. Neutron stars have the density $\rho \sim 10^{17}$–$10^{18}\,\mathrm{kg/m^3}$. Therefore, neutron stars fall to the horizon of black holes with a mass of $M > 20M_\odot$ without destruction. Of course, a collision with ultra-relativistic velocities of neutron stars in the vicinity of a massive black hole should be considered a very unlikely event. Estimates in [13] show that in the collision near the vicinity of the horizon of an extremal rotating black hole with a mass of $10^9 M_\odot$ at points with a radial coordinate $r_H + 7 \cdot 10^5$ km, the maximum collision energy in the center-of-mass system can reach 100 mc^2. In nucleon–nucleon collisions, this is the electro-weak unification energy. The masses of neutron stars range from one to three solar masses, and their radii are about 10–20 km. The gravitational radius of a black hole with a mass of 100 solar masses is approximately 300 km. Therefore, with such collision energy of two neutron stars, a black hole should form, and it will not be possible to obtain a substance in a state of an electro-weak phase transition outside the event horizon.

Thus, it is impossible to obtain the macroscopic quantities of a substance with an electro-weak phase transition with a lifetime significantly exceeding the Compton time for the electro-weak scale due to collisions in the vicinity of black holes.

5. The Influence of Spontaneous Symmetry Breaking on the Space–Time Metric

Let us consider a real scalar field with self-action [24]

$$V(\varphi) = -\frac{\mu^2}{2}\varphi^2 + \frac{\lambda^2}{4}\varphi^4 + \frac{\mu^2}{4\lambda^2}. \tag{26}$$

Here, $\mu = \tilde{\mu}c/\hbar$, $\tilde{\mu}$ is a mass parameter and λ is the dimensionless self-action constant. Stable equilibrium states of such a field are located at two points

$$\varphi = \pm\varphi_0, \quad \varphi_0 = \frac{\mu}{\lambda}. \tag{27}$$

The potential function (26) can be written as

$$V(\varphi) = \frac{\lambda^2}{4}\left(\varphi^2 - \varphi_0^2\right)^2. \tag{28}$$

Both lower states have zero energy, and the unstable equilibrium with $\varphi = 0$ has an energy density of

$$\varepsilon = \hbar c V(0) = \hbar c \frac{\mu^4}{4\lambda^2}. \tag{29}$$

Using the representation $\varphi = \varphi_0 + \chi$, one obtains

$$V(\chi) = \lambda^2 \varphi_0^2 \chi^2 + \lambda^2 \varphi_0 \chi^3 + \frac{\lambda^2}{4}\chi^4. \tag{30}$$

Thus, the mass of the χ field is $\sqrt{2}\lambda\varphi_0 = \sqrt{2}\mu$. In the case of the Higgs boson, $m_H = 125.3\,\mathrm{GeV}$, and we have

$$\varepsilon_H = \hbar c V(0) = \hbar c \frac{m_H^4 c^4}{16\hbar^4 \lambda^2} = \frac{1}{16\lambda^2}\frac{m_H c^2}{(l_C^H)^3}, \tag{31}$$

where $l_C^H = \hbar/(m_H c)$.

For the electro-weak interaction, quantum corrections lead to limitation [24]

$$\lambda \geq \alpha = \frac{e^2}{4\pi\varepsilon_0 \hbar c} \approx \frac{1}{137}, \tag{32}$$

where e is the elementary electric charge and ε_0 is the electric constant.

To estimate the inverse influence of the scalar field on the curvature of space–time, we use Einstein's equations

$$R_{ik} - \frac{1}{2}R g_{ik} + \Lambda g_{ik} = -8\pi \frac{G}{c^4}(T_{ik}^{(0)} + T_{ik}), \tag{33}$$

where Λ is the cosmological constant, $T_{ik}^{(0)}$ is the energy–momentum tensor of the background matter. The energy–momentum tensor for a constant scalar field with minimal coupling to curvature is [25]

$$T_{ik} = g_{ik}\hbar c V(\varphi) \tag{34}$$

and is similar to the contribution of an additional cosmological constant. Upon the appearance of a non-zero cosmological constant under spontaneous symmetry breaking, this was indicated in the work [26]. Phase transition in electro-weak interactions was discussed in cosmology by Kirzhnits and Linde [27,28], Weinberg [29] and others. Estimates of changes in the value of the cosmological constant during phase transitions in the early Universe were made in work [30].

If there is only a constant scalar field and the energy–momentum tensor of the background matter is equal to zero $T_{ik}^{(0)} = 0$, then the solution to Einstein's Equation (33) will be the de Sitter space–time. In de Sitter space, one has

$$R_{ik} = \frac{R}{4} g_{ik}, \tag{35}$$

and it follows from (33) that

$$R = 4\left(\Lambda + l_{\mathrm{Pl}}^2\, 8\pi V(\varphi)\right), \tag{36}$$

where l_{Pl} is the Planck length

$$l_{\rm Pl} = \sqrt{\frac{G\hbar}{c^3}} = 1.6162 \cdot 10^{-35}\,\mathrm{m}. \tag{37}$$

For the electro-weak case under $\varphi = 0$, from (31), we have

$$R = 4\left(\Lambda + \frac{\pi}{2\lambda^2}\frac{l_{\rm Pl}^2}{(l_C^H)^4}\right). \tag{38}$$

For the radius of curvature, we obtain

$$r \sim \frac{(l_C^H)^2}{l_{\rm Pl}} \sim 0.1\,\mathrm{m}. \tag{39}$$

This value is many orders of magnitude greater than the Compton wavelength of the particle and the size of the region in which the phase transition occurs. It should be expected that in order for special collisions with ultra-high energy to occur, in the volumes r^3, there must be a large number of particles falling onto the black hole. Their total mass will be much greater than the mass of the electro-weak scale. Thus, the inverse effect of energy density in the electro-weak phase transition in collisions on the space–time metric can be neglected.

6. Conclusions

An integral part of the standard model of particle physics is the mechanism of spontaneous symmetry breaking. The discovery at the Large Hadron Collider of the Higgs boson in 2012 makes us take seriously the possibility of a phase transition from one vacuum to another at high temperatures, as is the case in quantum nonrelativistic many-body theory, where the ground state plays the role of the vacuum. In our work [13], it was shown that in the processes of collisions of particles near the horizon of black holes, it is possible to achieve energies in the system of the center of mass of the order of the energy scale of the electro-weak phase transition.

In this article, we showed that the region of the phase transition in such collisions is microscopic. In the order of magnitude, the size of the region is equal to the Compton wavelength of the Higgs boson. Using formulas for black body radiation, we show that the lifetime of such region is of the order of the Compton time for the electro-weak phase transition scale.

During a phase transition, such as in the case of spontaneous symmetry breaking, the energy–momentum tensor corresponds to the emergence of an effective cosmological constant. It is shown that for phase transitions occurring during particle collisions, its influence on the space metric in the phase transition region can be neglected.

Note that despite the short time existence and microscopic volumes of a new phase of matter during an electro-weak phase transition in collisions in the vicinity of the black hole horizon, its very existence is of fundamental importance for the study of elementary particle physics in the ultra-high energy region, which is unattainable on Earth.

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