On Phase Transitions near Black Holes

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Received August 29, 2022; revised September 9, 2022; accepted September 15, 2022

It has been shown that temperatures near the horizon of rotating black holes can be about the phase transition temperature in the Standard Model with the Higgs boson. The distance from the horizon and gravitational and electromagnetic radiation emitted in collisions between particles have been numerically estimated.

DOI: 10.1134/S0021364022601907

1. INTRODUCTION

A quark–gluon plasma with a temperature of $4 \times 10^{12}$ K was obtained in collisions of gold ions in the relativistic heavy ion collider at the Brookhaven National Laboratory (United States) in 2010. The production of the quark–gluon plasma with a temperature of $5 \times 10^{12}$ K in collisions of lead ions at energies about several TeV per colliding nucleon at the Large Hadron Collider (LHC) was reported in 2012. However, the study of collisions of particles near rotating black holes\textsuperscript{[1–5]} shows that there is a natural supercollider where collision energies are much higher than those reached at the LHC. Consequently, the quark–gluon plasma produced in such collisions should have very high temperatures. It is of interest whether these temperatures can be as high as the phase transition temperature in the Standard Model with the Higgs boson. The distance from the horizon and gravitational and electromagnetic radiation emitted in collisions between particles have been numerically estimated.

The possible phase transitions in the early Universe lead to the known problem of the cosmological constant\textsuperscript{[9]}. Indeed, the energy density of vacuum corresponding to the mentioned phase transitions is many orders of magnitude higher than the value corresponding to the observed cosmological constant. To explain this, the fine tuning of the bare cosmological constant is necessary.

2. PHASE TRANSITIONS IN THE EARLY UNIVERSE

In the extrapolation of the standard cosmological model to times close to the Big Bang time, very high temperatures can theoretically be reached at which the following transitions between states of space matter can occur, usually called cosmological phase transitions.

(i) Transition between the quark–gluon plasma and hadrons at energies $E \approx 200$ MeV. The corresponding temperature $T = E/k_B \approx 10^{12}$ K, where $k_B \approx 1.380649 \times 10^{-23}$ J/K is the Boltzmann constant, can exist in the expanding Universe at a time of about $10^{-6}$ s after the Big Bang.

(ii) The electroweak phase transition at energies $E_W \approx 100$ GeV. The corresponding temperature $T_W \approx 10^{15}$ K can exist at a time of about $10^{12}$ s after the Big Bang.

(iii) The grand unification phase transition at energies $E_{GUT} \approx 10^{16}$ GeV. The corresponding temperature $T_{GUT} \approx 10^{29}$ K can hardly be reached in the early Universe\textsuperscript{[9]} according to models with an inflation stage in which the heating temperature is much lower than $T_{GUT}$. According to models with a radiation dominant stage, the temperature $T_{GUT}$ can be reached in the early Universe at times about $10^{-38}$ s.
3. HAWKING TEMPERATURE NEAR THE HORIZON OF A NONROTATING BLACK HOLE

High temperatures and regions of possible phase transitions can occur near the horizon of black holes owing to the thermal emission of black holes discovered by Hawking [10, 11].

The metric of a nonrotating black hole can be represented in the form

\[ ds^2 = \left(1 - \frac{r_g}{r}\right) c^2 dt^2 - \frac{dr^2}{1 - \frac{r_g}{r}} - r^2 d\Omega^2. \]  

(1)

Here, \( r_g = 2GM/c^2 \), where \( G \) is the gravitational constant and \( M \) is the mass of the black hole; \( c \) is the speed of light in vacuum; and \( d\Omega^2 = d\theta^2 + \sin^2 \theta d\phi^2 \). The Hawking temperature at infinite distance from the Schwarzschild black hole is [12]

\[ T_H = \frac{\hbar c^3}{8\pi k_B GM} = 6.169 \times 10^{-8} \frac{M_c}{M}, \]  

(2)

where \( \hbar \approx 1.05457 \times 10^{-34} \text{ J s} \) is the reduced Planck constant and \( M_c \) is the mass of the Sun.

As known [13],

\[ T\sqrt{g_{00}} = \text{const} \]  

(3)

in the gravitational field in thermodynamic equilibrium. Consequently, the Hawking temperature of the nonrotating black hole at points with the radial coordinate \( r \) is given by the expression

\[ T(r) = T_H \sqrt{1 - \frac{r_g}{r}} = \frac{\hbar c}{4\pi k_B r_g} \frac{r}{\sqrt{\Delta r}} \]  

(4)

\[ = 6.169 \times 10^{-8} \frac{M_c}{M} \sqrt{\frac{r}{\Delta r}}, \]

where \( \Delta r = r - r_g \). For radial coordinates \( r = r_g \) at which the temperature of the Hawking radiation is \( T \),

\[ \Delta r = \frac{\hbar^2 c^4}{32\pi^2 k_B^2 GM^2} \approx 1.12 \times 10^{-11} \frac{M_c}{M} \text{ (m)} \]  

(5)

Here, the temperature is given in kelvins and the distance is in meters. The substitution of the quark–gluon phase transition temperature and \( M = M_c \) into Eq. (5) gives \( \Delta r \approx 10^{-35} \text{ m} \), which is about the Planck length \( l_p = \sqrt{\hbar G/c^3} = 1.616 \times 10^{-35} \text{ m} \). This distance for black holes with masses larger than the mass of the Sun is even smaller.

It is noteworthy that the radial coordinate \( r \) is not identical to the physical distance, which can be defined only locally in the curved spacetime [14]. To illustrate the meaning of the radial coordinate, we present the radial equation of timelike geodesics in the Schwarzschild field (1) (see [15, Section 19]):

\[ \left(\frac{dr}{c d\tau}\right)^2 = \varepsilon^2 - \left(1 - \frac{r_g}{r}\right) \left(1 + \frac{L^2}{(mc^2)^2}\right), \]  

(6)

where \( \tau, \varepsilon = E/(mc^2) \), and \( L \) are the self-time, specific energy, and angular momentum of the moving particle with the mass \( m \). According to Eq. (6),

\[ d\tau = \frac{c}{\sqrt{\varepsilon^2 - 1 - \frac{r_g}{r} \left(1 + \frac{L^2}{(mc^2)^2}\right)}} \frac{dr}{c}. \]  

(7)

According to Eq. (7) at \( r \rightarrow r_g \), the residence time of the particle with a fixed angular momentum in the region \( dr/c \) is \( \approx dr/(c\varepsilon) \). The total energy of the particle incident from a region far from the event horizon cannot be much lower than \( mc^2 \), and, therefore, the residence time in the indicated region does not exceed \( dr/c \) in order of magnitude. Thus, the above estimates of \( \Delta r \) for the quark–gluon phase transition show that matter outside the event horizon can be in such state for a time about the Planck time, which is physically unobservable. Consequently, phase transitions, even quark–gluon ones, near the horizon of black holes at the Hawking temperature cannot be observed.

4. TEMPERATURE REACHED IN COLLISIONS NEAR THE HORIZON OF AN EXTREME ROTATING BLACK HOLE

The Kerr metric of a rotating black hole [16] in the Boyer–Lindquist coordinates [17] has the form

\[ ds^2 = \frac{\rho^2}{\Sigma^2} c^2 dt^2 - \frac{\sin^2 \theta}{\rho^2} (d\phi - \omega dt)^2 - \frac{\rho^2}{\Delta} dr^2 - \rho^2 d\theta^2, \]  

(8)

where

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

\[ \Sigma^2 = \left(r^2 + \frac{a^2}{c^2}\right) - \frac{a^2}{c^2} \sin^2 \theta \Delta, \quad \omega = \frac{2Gma}{\Sigma^2 c^2}, \]  

(9)

\[ M \text{ and } aM \text{ are the mass and angular momentum of the black hole, respectively. We accept that } 0 \leq a \leq GM/c. \text{ The event horizon of the Kerr black hole has the radial coordinate } \]  

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

(10)

According to [1], the squared energy of collision of two particles with the mass \( m \) with the angular momenta \( L_q \) and \( L_z \) in the center-of-mass system, which are nonrelativistic at infinity and are freely inci-
The energy of collision of two particles with the angular momenta \( l_1 = 4 \) and \( l_2 = -4 \) near the event horizon of the nonrotating \( (A = 0) \) black hole is \( 2\sqrt{5}mc^2 \). The maximum achievable energy of collision in the center-of-mass system increases with the velocity of rotation of the black hole [3].

The particle with the maximum possible angular momentum \( l_i = 2 \) (critical particle) freely incident on the event horizon for an infinitely long self-time. The energy of collision of this critical particle with another particle with the angular momentum \( l_i \) in the interval of \((-1, 2)\) near the event horizon can be unlimitedly high (Banados–Silk–West resonance [1]):

\[
\frac{E_{c.m}^2}{m^2c^4} = \frac{4}{x(x-1)} \left(1 - l_i + x^2 - \sqrt{(1- l_i^2) - l_i^2x/2 + x^2} \right).
\]

The substitution of this energy into the formula \( T = (E_{c.m} - 2mc^2)/k_B \) for the temperature yields

\[
T = \frac{2mc^2}{k_B} \times \left[ \frac{1 - l_i + x^2 - \sqrt{(1- l_i^2) - l_i^2x/2 + x^2}}{x(x-1)} - 1 \right]
\]

\[
= 1.083 \frac{mc^2}{k_B} \sqrt{\frac{2 - l_i}{x-1}} \quad x \to 1.
\]

This temperature is reached at the distance

\[
r - r_H = 1.17(2 - l_i)r_H \left(\frac{mc^2}{k_Bc}\right)^2, \quad r \to r_H.
\]

For the mass \( m \) about the proton mass and \( l_i = 0 \), the electroweak temperature near the extreme rotating black hole is reached at the distance

\[
r - r_H = 2 \times 10^{-4} r_H.
\]

This distance for black holes with the mass of the Sun is tens of centimeters and, in contrast to the Hawking temperature, is acceptable for the occurrence of a phase transition.

To estimate the corresponding temperature far from the black hole, we use the equation for the time component of the timelike geodesic [15]:

\[
\rho \frac{d\tau}{dt} = \frac{1}{\Delta} \left( \Sigma^2 \varepsilon - 2 \left( \frac{GM}{c^2} \right)^2 \right) \tau \Delta.
\]

For the particle with \( l < 2\varepsilon r_Hc/a \) (noncritical particle) freely incident near the horizon of the rotating black hole, the dilation of the self-time \( \tau \) compared to the time at infinity \( t \) is

\[
\frac{d\tau}{dt} \sim \frac{2\rho(2\rho - l)}{(x-1)^2(1 + \cos^2\theta)}, \quad x \to 1.
\]

The dilation of the self-time for the critical particle with \( l = l_H = 2\varepsilon r_Hc/a \) is

\[
\frac{d\tau}{dt} \sim \frac{4\rho}{(x-1)(1 + \cos^2\theta)}, \quad x \to 1.
\]

Constraints on the possible specific energy \( \varepsilon \) and the projection of the angular momentum \( l \) on the axis of rotation of the black hole at a given radial coordinate \( r \) can be obtained from equations for the radial and angular components of geodesics. We consider only the motion in the equatorial plane with \( \theta = \pi/2 \). In this case, the equation for the radial component of a geodesic has the form [15]:

\[
\frac{\rho^2}{c} \frac{d\tau}{dt} = \pm \sqrt{R},
\]

\[
R = \Sigma^2 \varepsilon^2 - l^2 \left( \frac{GM}{c^2} \right)^2 \times \left( 4\rho \frac{a}{c} + l \left( r - 2 \frac{GM}{c^2} \right) \right) - r^2 \Delta.
\]

For the critical particle, the condition \( R \geq 0 \) gives the constraint \( \varepsilon \geq \sqrt{x/(x+2)} \); hence, the time dilation is

\[
\frac{d\tau}{dt} \geq \frac{4}{\sqrt{3(x-1)}}, \quad x \to 1.
\]

The specific energy of noncritical particles incident on the extreme black hole can be low. In particular, the condition \( R \geq 0 \) in the case \( l = 0 \) yields \( \varepsilon \geq \varepsilon_{\text{min}} \), where

\[
\varepsilon_{\text{min}} = \frac{(x-1)\sqrt{x}}{\sqrt{x^3 + x + 2}} \sim \frac{x - 1}{2}, \quad x \to 1.
\]
In the limiting case \( \varepsilon = \varepsilon_{\text{min}} \), according to Eq. (18), the time dilation is

\[
\frac{dt}{d\tau} = \frac{\sqrt{x^2 + 1} + 2/x}{x - 1} \sim \frac{2}{x - 1}, \quad x \to 1.
\] (25)

The collision of particles near the event horizon of the black hole with the production of the quark–gluon plasma can be accompanied by the production of particles with various specific energies and momenta and the subsequent emission of photons at their collisions. According to Eqs. (19), (23), and (25), the energy of a collision occurring at the point approaching the event horizon of the black hole \( x \to 1 \), as well as the local temperature of produced particles, can unlimitedly increase, but the temperature observed far from the black hole tends to zero owing to the redshift.

Combining Eqs. (15) and (25), we obtain the following approximate expression for the temperature observed far from the black hole as a function of the radial coordinate \( x \) of the collision point:

\[
T(x) \approx \frac{mc^2}{k_B} \frac{2(x - 1)}{\sqrt{x^2 + x + 2/x}} \times \left[ \frac{1 - l_i + x^2 - \sqrt{(1 - l_i)^2 - l_i^2x/2 + x^2 + 2}}{x(x - 1)} \right].
\] (26)

The maximum \( T(x) \) value at \( l_i = -4 \) is \( \approx 0.534mc^2/k_B \) and is reached at \( x = 1.744 \).

The corresponding expression in the case of the collision of particles with \( l_i = -4, \ l_2 = 4 \), and \( E_{1,2} = mc^2 \) incident on the Schwarzschild black hole has the form

\[
T(x) = \frac{mc^2}{k_B} \frac{32}{x(x + \sqrt{x^2 + 16})} \sqrt{\frac{x - 2}{x}}.
\] (27)

In this case, the maximum \( T(x) \) value is \( \approx 0.803mc^2/k_B \) and is reached at \( x = 2.645 \).

Thus, the temperature observed far from high-energy collisions near the event horizon of the black hole rotating at any velocity is in order of magnitude no more than several tenths of \( mc^2/k_B \). The local temperature near the collision point near the extreme rotating black hole can be unlimitedly high. Because of the conservation laws, the total energy emitted from the black hole cannot be higher than the sum of the energies of the colliding particles \( 2mc^2 \) for the collision of particles with the same mass \( m \), which are nonrelativistic at infinity, if the Penrose effect for rotating black holes is disregarded.

It is remarkable that the extreme rotating black hole is no longer extreme after the incidence of any particle on it [18]. Consequently, extreme black holes can hardly exist in nature. The extreme angular momentum of the black hole achievable upon the accretion of matter on it is estimated in [19] as \( A = 0.998 \). At this angular momentum of the black hole, the maximum energy of a collision in the center-of-mass system of freely incident particles nonrelativistic at infinity is only \( E_{\text{c.m.}} \approx 19mc^2 \).

5. MULTIPLE COLLISIONS NEAR THE HORIZON OF NONEXTREME ROTATING BLACK HOLES

As shown in [2, 3], a superhigh 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. The possible collision energy in the center-of-mass system of the particle with the angular momentum \( l \) and the other particle with the same mass \( m \) can be estimated as [2]

\[
E_{\text{c.m.}}(r) = \frac{mc^2}{\sqrt{\delta}} \frac{2(l_H - l)}{\sqrt{1 - \sqrt{1 - A^2}}}.
\] (28)

where \( \delta \) is the parameter characterizing the region of radial coordinates where the high-energy collision is possible:

\[
x \leq x_H + \frac{\delta^2(1 - \sqrt{1 - A^2})^2}{4x_H \sqrt{1 - A^2}}.
\] (29)

According to Eqs. (28) and (29), the distance at which the collision energy \( E_{\text{c.m.}} \) can be reached is given by the expression

\[
x - x_H = \frac{(l_H - l)^2}{x_H \sqrt{1 - A^2}} \left( \frac{mc^2}{E_{\text{c.m.}}} \right)^4.
\] (30)

In particular, for rotating black holes with \( A \approx 1 \) at \( l = 0 \),

\[
r - r_H = \frac{2.8r_H}{\sqrt{1 - A}} \left( \frac{mc^2}{E_{\text{c.m.}}} \right)^4.
\] (31)

At the limiting value \( A = 0.998 \) [19], the electroweak temperature in order of magnitude is possible at \( r - r_H \approx 6 \times 10^{-3}r_H \). This distance is about meters for black holes with the mass of the Sun and can be thousands of kilometers for supermassive black holes.

Thus, the quark–gluon and even electroweak phase transition temperatures can be reached in multiple collisions of particles near rotating black holes.
6. Collisions of Macroscopic Bodies

Collisions of macroscopic bodies near the event horizon of black holes are possible only when such bodies reach the horizon region rather than being destroyed by tidal forces of the gravitational field. Many processes of destruction of macroscopic objects (stars) by tidal forces near supermassive black holes were observed in the SRG/erosita space experiment [20]. Tidal forces near the horizon decrease with an increase in the mass of the black hole. We present below some estimates of the mass of black holes at which macroscopic objects near the event horizon are not destroyed (see also [21, p. 772]). We consider only the nonrotating black hole and radial tidal forces. According to the equation for the deviation of geodesics in the coordinates associated with the center of mass of the incident body (see [22, Eq. (32.24b)]),

\[ \frac{D_{\xi}^2}{d \tau^2} = \frac{2GM_{\xi}}{r^3}, \]

where \( \xi \) is the corresponding radial coordinate. For an estimate, we assume that a star or an incident planet is destroyed if tidal forces acting on 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 incident body. The condition for the incidence of a uniform ball with the density \( \rho \) and radius \( R \) to the horizon has the form

\[ \frac{2GM}{r_g^3} R < \frac{G^4 \pi \rho R^3}{3R^2}. \]

Therefore, the star or planet bound by the gravitational forces is not destroyed upon incidence to the event horizon if the mass of the black hole satisfies the inequality

\[ M > \frac{c^4}{4G^{3/2}} \sqrt{\frac{3}{\pi \rho}} \frac{M}{M_*} > 1.9 \times 10^8 \sqrt{\frac{\rho_w}{\rho}}. \]

where \( \rho_w = 10^3 \text{ kg/m}^3 \). Thus, a collision of stars or planets with the relative velocity close to the speed of light in vacuum is possible only near the event horizon of supermassive black holes with the mass exceeding \( 10^8 \) of the mass of the Sun. Such black holes exist at the centers of many galaxies. The collision of compact objects with star masses near supermassive black holes was considered in [23]. The authors of [24] showed that large clouds of comets, asteroids, and stones should be formed around supermassive black holes.

Let us estimate the mass of the black hole allowing incidence of an undestroyed stone with a size of several centimeters/meters. The incident stone is destroyed if tidal forces exceed its tensile strength for the radial direction or its compressive strength for the polar and azimuthal directions. We consider only the case of breaking in the radial direction and take the tensile strength of titanium alloy \( \sigma = 10^9 \) Pa. The necessary condition for the body with the characteristic size \( d \) and density \( \rho \) not to be destroyed has the form

\[ \frac{GM}{4r_g^3} \rho d^2 < \sigma \]

(see the derivation of Eq. (32.25a) in [22]). Substituting the expression \( r_g = 2GM/c^2 \) and the iron density \( \rho = 7.87 \times 10^3 \) kg/m\(^3\), we obtain

\[ M > \frac{c^4 d}{4G^{3/2}} \sqrt{\frac{\rho}{2\sigma}} \frac{M}{M_*} > 10^2 d, \]

where \( d \) is the distance in meters. Thus, collisions of 0.1-m stones with velocities close to the speed of light in vacuum are possible near the horizon of black holes with stellar masses.

The collision energy per each particle in the body rather than the total collision energy should be high to obtain high temperatures. The multiple collision mechanism is obviously inapplicable for macroscopic bodies because any relativistic collision inevitably destroys bodies. Consequently, collisions with center-of-mass energies much higher than \( mc^2 \) could be possible only near the horizon of extreme rotating black holes.

The process of ultrarelativistic collisions of macroscopic bodies expectedly has a complex structure. Such processes have not yet been observed in nature. Although the velocities of colliding neutron stars or merging black holes that are detected by gravitational radiation bursts are close to the speed of light in vacuum, the corresponding relativistic factor \( \gamma = 1/\sqrt{1 - v^2/c^2} \) only insignificantly exceeds unity.

The kinetic energy of each particle in colliding ultrarelativistic macroscopic bodies is much higher than the energy of electromagnetic bonds between these particles in the bodies. Therefore, such bodies in the model of ultrarelativistic collisions should be treated as colliding particle clouds. If the size of a cloud satisfies the inequality \( \sigma d \ll 1 \), where \( \sigma \) is the effective scattering cross section and \( n \) is the number density of particles, only an insignificant number of particles in the cloud are involved in the reaction and the colliding clouds pass through each other without destruction. For an estimate, we take the area of the cross section of a nucleus \( \sigma = \pi r_n^2 \) with the radius \( r_n \approx 10^{-15} \) m and the particle density of iron atoms \( n = 8.4 \times 10^{28} \) m\(^{-3}\). Then, colliding solid bodies with the characteristic sizes \( d \geq 30 \) m will be destroyed. In this case, the collision of macroscopic bodies will lead to ultrarelativistic pair collisions of nucleons constituting the bodies.

Electromagnetic and gravitational radiation bursts should be expected from ultrarelativistic collisions of macroscopic bodies. We assume that the collision of macroscopic bodies leads to pair collisions of nucleons constituting these bodies and use estimates obtained in

\( \text{JETP LETTERS} \quad \text{Vol. 116} \quad \text{No. 8} \quad 2022 \)
Gravitational radiation from pair collisions of particles with the mass $m$ can be estimated by Eq. (10.4.23) from [26]. This radiation will be suppressed by a factor of $n^2/M_p^2$ (see Eq. (11) in [25]), where $M_p$ is the Planck mass. The energy of electromagnetic radiation in order of magnitude is (see [25, Eq. (17)])

$$E_{em} = \frac{e^2}{\hbar c} E_{c.m.},$$

where $e$ is the elementary charge. The total intensity of electromagnetic radiation from the collision of macroscopic bodies will obviously be significant and can be observed even taking into account the redshift near the horizon. In particular, the maximum center-of-mass energy of the collision at points with a radial coordinate of $r_{H} + 7 \times 10^5$ km near the horizon of the extreme rotating black hole with the mass $10^9 M_\odot$ can reach $100mc^2$. In nucleon–nucleon collisions, this is the electroweak unification energy. The maximum total center-of-mass energy of the collision of two iron asteroids with a diameter of 1 km with the corresponding $\gamma$ factor is about $3 \times 10^{31}$ J. The center-of-mass energy of electromagnetic radiation near the collision point can be $2 \times 10^{29}$ J at a power of about $10^{34}$ W. Far from the black hole, an electromagnetic burst with an energy of about $10^{25}$ J at a power of about $10^{28}$ W can be observed owing to the redshift and time dilation in the gravitational field.

7. CONCLUSIONS

The Standard Model predicts the existence of the Higgs particle according to the Higgs model based on the Goldstone model. As known [27], the Goldstone model is similar to the theory of superfluidity and superconductivity and implies the existence of two vacua, symmetric and asymmetric. The experimental discovery of the Higgs boson at the LHC allows the possibility of a phase transition from one vacuum to the other at high temperatures, as in the nonrelativistic quantum theory of many bodies, where the ground state serves as vacuum. It has been shown in this work that such a phase transition in the Standard Model is possible in multiple collisions near the horizon of nonextremal rotating black holes (in their ergosphere). This phase transition is also possible in collisions of two macroscopic bodies near extremal black holes in the presence of the Banados–Silk–West resonance [1]. The emission of gravitational and electromagnetic waves accompanying the phase transition has also been discussed. Gravitational radiation in these collisions is insignificant, but electromagnetic radiation is fairly intense and the corresponding burst can be observed on the Earth.

FUNDING

This work was supported by the Russian Science Foundation (project no. 22-22-00112).

CONFLICT OF INTEREST

The authors declare that they have no conflicts of interest.

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