frontiers | Frontiers in Physics

TYPE Original Research
PUBLISHED 13 June 2025
DOI 10.3389/fphy.2025.1599363

Check for updates

OPEN ACCESS

EDITED BY
Marco Muccino,
National Laboratory of Frascati (INFN), Italy

REVIEWED BY
Orlando Luongo,
University of Camerino, Italy
Maryam Azizinia,
University of Camerino, Italy

*CORRESPONDENCE
Johannes Böhm-Mäder,
j.boehm@universe20.ch

RECEIVED 24 March 2025
ACCEPTED 05 May 2025
PUBLISHED 13 June 2025

CITATION
Böhm-Mäder J (2025) Universe 2.0: Black holes? No. Black stars!
Front. Phys. 13:1599363.
doi: 10.3389/fphy.2025.1599363

Universe 2.0: Black holes? No. Black stars!

Johannes Böhm-Mäder*

Retired, Bubikon, Switzerland

Strong analytical evidence reveals that *every* solution of the Schwarzschild-Kerr metric, when continuously solving the metric across the event horizon, violates the conditions of the symmetry group, inevitably, which renders the long-term vision of black holes untenable. General relativity does not support a passable horizon hosting an infinitely concentrated mass in vacuum. The hitherto notion of black holes breaks down altogether. A new Schwarzschild solution, integrated within a massive shell, and a revaluation of the gravitational collapse yield a new vision of black stars of fluid, formed by a continuous progression of stacked horizons from the center to the surface. Thus, the current theory fundamentally differs from attempts to avoid horizons, such as mimickers, gravastars, boson stars, or gravitational metamaterials. The exterior view of Schwarzschild-Kerr black holes is not affected at all, but the inner view is obsolete. The information paradox and the conflict with the Pauli exclusion principle are solved and the Newtonian shell theorem is fully applicable in the relativistic context.

KEYWORDS

black hole, black star, Birkhoff theorem, event horizon, general relativity, information paradox, Pauli exclusion principle, Schwarzschild metric

1 Introduction

Only a few months after Einstein had released his equations of general relativity, Schwarzschild [1] published his spherically symmetric solution, which is irregular at the event horizon. About 5 years thereafter Painlevé [2] presented the first globally regular coordinates. In the same year Jebsen [3] provided the proof of the Birkhoff theorem. All this was more than a century ago. Various globally regular solutions followed until Kruskal [4] and Szekeres [5] independently found the maximal extension of the spherically symmetric metric. Hereupon, the scientific community agreed that the patches of an irregular manifold may be connected by singular transformations. For decades the globally regular solutions governed our notion of black holes as singularities that infinitely concentrate all structured energy passing the event horizon.

In more recent studies a number of authors, in view of the inconsistencies of the long-term notion, provided regular metrics of black holes by avoiding horizons. Lemos and Zaslavskii [6] proposed black hole mimickers that "would look almost like black holes but have ho horizon". Mimickers compared to black holes are stars in a critical state just before the gravitational breakdown. According to Casadio et al. [7], anisotropic pressures within the stellar structure might yield anti-gravitational effects and thus prevent the collapse. Yet, many questions regarding stability remain open. The mimickers might eventually collapse to black holes and thus lose their material integrity.

Mazur and Mottola [8] by a thin phase boundary separate an exterior Schwarzschild metric from an interior de Sitter region of $p = -\rho$ gravitational vacuum condensate. The phase boundary replaces the horizon. Even though it has no horizon, the exterior view of these gravastars exactly resembles Schwarzschild black holes. The interior de Sitter metric requires a negative pressure stabilizing the boundary, a notion similar to the hypothetical concept of dark energy, meant to explain the accelerated expansion of the universe. The gravastars also depend on a quantum gravitational vacuum phase transition before the classical event horizon can form.

Boson stars, first proposed by Tkachev [9], build on scalar fields and require postulating a stable type of massive bosons with self-repulsive interaction. The concept of boson stars received a boost by the discovery of the Higgs boson, but the latter is highly instable and no other massive bosons have been discovered yet. A wide variety of possible types of boson stars have been theoretically discussed (Liebling and Palenzuela, [18]), but to date, no unequivocal observational evidence has been found. A compact boson star by its gravitation would bend light and create an empty space-time region resembling the shadow of a black hole's event horizon. Thus, boson stars might resemble black holes, even though they do not have a horizon.

In a most resent preprint, Luongo [10] avoids the singularity by speculating on the space-time behaving like a gravitational metamaterial, relying on the hypothetical metamaterials proposed by Yang et al. [11]. In Luongo's concept a black hole's gravitational field is modeled as an exotic medium with a negative refractive index. The metamaterials may also exhibit a particle-like behavior that might contribute to dark matter.

The present study follows a completely different ansatz compared to mimickers, gravastars, boson stars, or metamaterials. It bases on the Schwarzschild or Kerr metric alone and requires no additional conditions or assumptions. We accept the gravitational singularities and implement them consistently. The analytical evidence presented in what follows naturally excludes black holes in the sense of an infinitely concentrated mass in a bubble of vacuum. A new conception results by evaluating the Schwarzschild solution within a concentric shell of mass instead of vacuum to infinity and by reanalyzing the gravitational collapse. The result can be extended to the Kerr metric. We end up in a vision of a black star, that is, of a continuous progression of stacked horizons from the center to the surface, forming a stratified body of fluid. The exterior view is not affected at all, but the long-term inner view as black holes is obsolete. The black stars are not an alternate view, but the unavoidable *replacement* of Schwarzschild-Kerr black holes.

2 Methods

In Section 3.1 we evaluate the conditions of the symmetry group. A scarcely considered reservation reveals an inherent discontinuity of the spherically symmetric metric, at the event horizon, where the Birkhoff theorem does not apply. The globally regular solutions are continuous, but at the cost of violating the symmetry conditions on the inner orbits. In Section 3.2, an inspection of the proper time of a freely falling observer descriptively confirms that in a physical space-time the event horizon cannot be passed.

A different vision of black holes arises in Section 4 by evaluating the Schwarzschild solution within a concentric spherical shell of mass instead of the boundary condition at infinity. The result is equal to the original Schwarzschild solution but extends the validity of the relativistic Newtonian shell theorem. The findings of Sections 3 and 4 allow reconsidering the gravitational collapse in Section 5. The discussion in Section 6 extends the new vision of black stars to the Kerr solution.

3 Revision of the spherically symmetric solutions

In the current section we revise the spherically symmetric vacuum solutions of general relativity. Section 3.1 follows up a reservation of the metric, which was only rarely considered in the literature. Section 3.2 complements the results by inspecting a freely falling observer. The concluding summary in Section 3.3 consolidates the arguments that challenge the hitherto notion of Schwarzschild-Kerr black holes.

3.1 Conditions of the symmetry group

According to the Birkhoff theorem, all the spherically symmetric vacuum solutions of general relativity form a class of solutions of one and the same Schwarzschild metric. Its prove proceeds from a spherically symmetric ansatz with a general, time dependent line element, for instance,

$$ds^2 = -A(t,r)\,dt^2 + 2B(t,r)\,dt\ dr + C(t,r)\,dr^2 + D(t,r)\,d\Omega^2\ , \quad (1)$$

where

$$r^2 d\Omega^2 = r^2\left(d\vartheta^2 + \sin^2\vartheta\ d\varphi^2\right) \quad (2)$$

denotes the surface area element of a concentric 2-sphere with a radius r. The metric Equation 1 is invariant upon a transformation

$$(t,r) \to \left(\tilde{t}(t,r), \tilde{r}(t,r)\right). \quad (3)$$

If we translate the coordinate t by a function $\psi(t,r)$ such that

$$\frac{d\psi(t,r)}{dr} = -\frac{B(t,r)}{A(t,r)} \quad (4)$$

and choose the transformed $\tilde{D} = r^2$, Equation 1 takes the form

$$ds^2 = -\tilde{A}(t,r)\,dt^2 + \tilde{C}(t,r)\,dr^2 + r^2\,d\Omega^2. \quad (5)$$

The translation by $\psi(t,r)$ preserves $A(t,r)$, that is, $\tilde{A} = A$. Equation 5 represents the standard form of the spherically symmetric metric, which serves to prove the Birkhoff theorem. For the translation by $\psi(t,r)$ we tentatively should consider the ancillary condition $A(t,r) \neq 0$, where Equation 4 is singular or undefined. In the resulting Schwarzschild metric, $\tilde{A} = 1 - 2m/r$ with the central point mass m (in units with the light speed $c = 1$ and the gravitational constant $G = 1$). Thus, $A(t,r) \neq 0$ corresponds to $r \neq 2m$.

Regarding a transformation as in Equation 3, Plebański and Krasiński, [12] (P&K, pp. 168–171) discuss a reservation. In the spherically symmetric manifold, the coordinates ϑ and φ from Equation 2 define a point on a 2-sphere with constant values of the arbitrary coordinates t and r. The function $D(t,r)$ in Equation 1 is a scalar on the orbits of the symmetry group. In the resulting metric (Equation 5) the gradient of the radius of the 2-sphere is space-like, $g^{\mu\nu}\partial_\mu r \partial_\nu r > 0$. Because $D(t,r)$ transforms as a scalar, this statement is

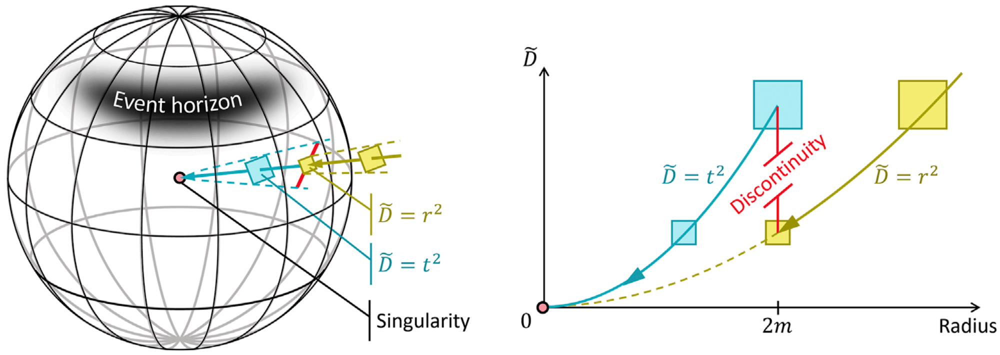

FIGURE 1
Schematic visualization of the inherent discontinuity of the spherically symmetric metric, due to the scalar $\tilde{D}(r > 2m) = r^2$ (yellow, P&K case I) necessarily differing from $\tilde{D}(r < 2m) = t^2$ (blue, P&K case II). The discontinuity at $r = 2m$ (P&K case III) is marked in red.

invariant upon a transformation like Equation 3. Thus, the gradient must have been space-like in the original metric (Equation 1), $g^{\mu\nu}\partial_\mu D(t,r)\partial_\nu D(t,r) > 0$.

P&K distinguish four cases, three of which are relevant for the present study. Case I relates to the exterior of the Schwarzschild metric, $r > 2m$ ($\tilde{A} > 0$), where the gradient of $D(t,r)$ is a space-like vector, which allows for $\tilde{D} = r^2$, as in Equation 5. The interior, $r < 2m$, indicates case II, for which P&K show the gradient of $D(t,r)$ being a time-like vector, $g^{\mu\nu}\partial_\mu D(t,r)\partial_\nu D(t,r) < 0$. This corresponds to t and r interchanging their roles in the metric and excludes $\tilde{D} = r^2$, but instead indicates $\tilde{D} = t^2$.

In the transition from case I to case II, the time t and the radius r interchange their nature, on a path, which in the angular directions is measured by a scalar immediately switching from $\tilde{D} = r^2$ to $\tilde{D} = t^2$. The interchange may be accepted as a reinterpretation of the coordinates or a change in the space-time topology. The switch of the scalar $\tilde{D}$, however, represents a discontinuity between adjacent orbits of the symmetry group. Figure 1 schematically visualizes the behavior of $\tilde{D}$.

For $\tilde{A}(t,r) = 0$, that is, for $r = 2m$, the gradient of $D(t,r)$ is a nontrivial null vector, $g^{\mu\nu}\partial_\mu D(t,r)\partial_\nu D(t,r) = 0$. This implies P&K's case III, which allows for $\tilde{D} = r^2$ but excludes $\tilde{B} = 0$, because the determinant $\tilde{A}\tilde{D} - \tilde{B}^2$ would vanish. Consequently, the metric in Equation 1 at $r = 2m$ cannot be transformed to the diagonal metric in Equation 5. This confirms the consideration regarding Equation 4. The standard form of the spherically symmetric metric and thus the Birkhoff theorem are restricted to $r \neq 2m$.

As an example of the relation between the solutions of the Schwarzschild metric, we consider P&K's sequence of transformations in their section 14.9, leading to the Kruskal-Szekeres coordinates. Similar to the Eddington-Finkelstein solution a translation of the time by $r^* = r + 2m \ln|r/2m - 1|$ is involved, which removes the singularity of $g_{rr} = (1 - 2m/r)^{-1}$. Figure 2 depicts the function $\ln|r/2m - 1|$.

All the transformations leading to globally regular coordinates of the Schwarzschild metric use a similar, singular function, Painlevé-Gullstrand, for example, $\ln|\sqrt{r/2m} - 1|$. Considerable effort is spent on arguments that the transformations are globally single-valued and invertible. However, a coordinate transformation involving a function as depicted by Figure 2 should be considered nonisomorphic, locally at $r = 2m$, where it is singular and thus discontinuous.

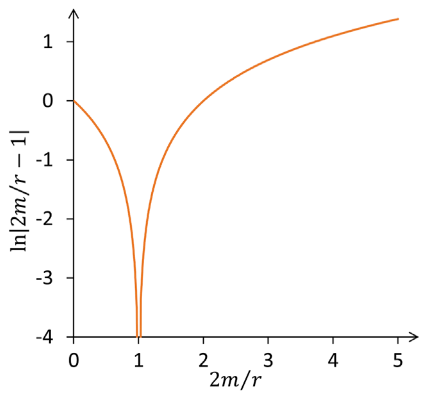

FIGURE 2
Graph of the transformation function $\ln|r/2m - 1|$ (orange line), with the singularity at the radius $r = 2m$.

Thereafter, P&K (p. 192) introduce an arbitrary constant a to be assigned later. Their equation 14.100 in their transformed coordinates $p(t,r)$ and $q(t,r)$ reads

$$ds^2 = -\left(1 - \frac{2m}{r}\right)a^2 e^{-2r/a}\left(\frac{r}{2m} - 1\right)^{-4m/a} \mathrm{d}p\ \mathrm{d}q + r^2 \mathrm{d}\Omega^2. \quad (6)$$

Then, they choose $a = 4m$, such that the factor causing the singularity cancels out. At $r = 2m$, this results in a division 0/0. The resulting metric may be defined at $r = 2m$, the transformation, however, is not. The gradient of the scalar $D(t,r)$ in the general metric in Equation 1 changes from space-like to time-like. In the transition to $r < 2m$, Equation 6 and eventually the Kruskal-Szekeres coordinates would have to switch the angular components from $r(p,q)^2\mathrm{d}\Omega^2$ to $t(p,q)^2\mathrm{d}\Omega^2$.

We conclude that we should divide the Birkhoff class of solutions in two subclasses, the first one derived from the Newtonian limit, the other related to it by transformations involving a singular function, as depicted by Figure 2. They all are valid solutions of Einstein's field equations, at least for $r > 2m$. Yet, locally at $r = 2m$ the two subclasses are related by nonisomorphic transformations and thus do not represent the same manifold. Assuming the region $r < 2m$

as well-defined, all the globally regular coordinates would have to switch the angular components of the line element from $r^2 d\Omega^2$ to $t^2 d\Omega^2$. Nobody ever appears to do so. Otherwise, the metric would reveal its inherent discontinuity. The inner solutions would even conflict with Einstein's vacuum equations.

3.2 Vertical free fall in the Schwarzschild metric

The original Schwarzschild solution is the only member of the Birkhoff class approaching the Newtonian limit. Thus, it is the suitable choice for imbedding a black hole into its surroundings. The proper time of a stationary observer at great distance in the surrounding universe corresponds to the metric's coordinate time.

We consider in the Schwarzschild metric the proper time τ of an observer in vertical free fall ($d\vartheta/d\tau = d\varphi/d\tau = 0$). The proper time span required to cross $r = 2m$ and to reach $r = 0$ is integrable and finite. While the falling observer approaches $r = 2m$, the coordinate time t diverges. Even though the integrated τ is finite, the falling observer from the viewpoint of a stationary observer at $r \gg 2m$ appears never to reach $r = 2m$.

The irregularity of the Schwarzschild solution is rated spurious, the singularity a mere coordinate singularity. Thus, it was generally concluded that there must be something wrong with the diverging coordinate time t. However, the proper time of a stationary observer at great distance in the free space-time is a good measure of the overall evolution of the universe. Whether we assume the latter closed, flat, or open, in the eternity the black hole would no longer exist and no distant observer would be left to witness the event.

Reconsidering the determination of the proper timespan for the falling observer's path, we find τ integrated over the coordinate radius r. The proper radius of the falling observer diverges at $r = 2m$, but the time component g_{tt} of the Schwarzschild coordinates vanishes inversely proportional to g_{rr}. Thus τ may asymptotically reach the corresponding limit τ_S.

In the region $r < 2m$, however, the proper time of any distant stationary observer in the universe with respect to the falling observer is undefined. The issue in integrating τ beyond $r = 2m$ lies in t and r interchanging their roles in the metric. In P&K's case II discussed in Section 3.1 above, the gradient of the 2-sphere's radius is a time-like vector. Thus, the proper time at $r < 2m$ would have to be defined with respect to the time-like coordinate r and integrated over the space-like coordinate t.

We conclude that the common integration of the finite proper time of a falling observer to reach $r = 0$ is invalid. Moreover, the analysis above provides an indication that at $r < 2m$ time and radius do not interchange their roles. The stationary time measure in the surrounding universe with respect to a falling observer is undefined, in descriptive terms, beyond infinity. There is no continuous path to $r < 2m$.

3.3 Summary of the revision

The findings regarding the Schwarzschild metric consolidate to the following statements.

1. The standard form of the spherically symmetric metric (Equation 5) due to the conditions of the symmetry group is subject to the ancillary constraint $r \neq 2m$.
2. The Birkhoff theorem, depending on the standard form, is subject to the same constraint $r \neq 2m$.
3. The irregular Schwarzschild coordinates exclusively approach the Newtonian limit. Thus, they describe a physical space-time, at least for $r > 2m$.
4. The globally regular coordinates are valid solutions of Einstein's field equations but are related to the Schwarzschild solution by transformations that are discontinuous and thus nonisomorphic at $r = 2m$.
5. Because of the statements 2 and 4, the globally regular solutions may be considered a separate subclass of the Birkhoff theorem, at $r = 2m$ belonging to a different manifold that might not represent a physical space-time.
6. For the proper times t of a stationary observer at $r \gg 2m$ and τ of a freely falling observer, the 2-sphere $r = 2m$ represents the asymptotical boundary $(t, \tau) \to (\infty, \tau_S)$, whereas $\tau > \tau_S$ is undefined (imaginary).
7. Geodesics extending beyond the event horizon interpolate over the discontinuity of the physical space-time that represents the eternity beyond the end of the universe, at $r < 2m$ violating the conditions of the symmetry group.

By very similar arguments it could be shown that the Kerr [13] metric for rotating black holes should be subject to the ancillary constraint $r^2 - 2mr - a^2 \neq 0$, where a denotes the Kerr parameter related to the angular momentum.

These findings challenge the current notion of Schwarzschild-Kerr black holes. If the event horizon is impassable from the outside, a black hole cannot be a gravitational annihilator engulfing in its central singularity all the energy approaching the point of no return.

It is hard to figure why nobody noted the inherent discontinuity of the symmetry group or its conditions being violated, not even P&K, who revealed the details. We can only argue as follows. The community needed a concept beyond the event horizon and Schwarzschild's original solution did not permit an extension to $r < 2m$. His inner solution [14] with the diverging pressure was available from the very beginning. The material collapse into the central singularity appeared as the most natural behavior. With no alternative at hand, there was no incentive to challenge the notion that was solidly established for decades. Only lately, as reported in the Introduction above, several authors searched for alternate metrics of black holes, avoiding the singularity. Yet, based on the irregular Schwarzschild solution, there would be a loophole allowing for a different view of the Schwarzschild metric, as elaborated in Section 4.

4 Exterior Schwarzschild solution with a new boundary

In the current section, we confront the finding from the preceding Section 3 that the interior of a black hole is inaccessible from the outside. We trace the Schwarzschild solution, replacing the boundary condition.

Schwarzschild [1] solved Einstein's equations for the boundary condition at $r \to \infty$ in vacuum. Very similar conditions apply within a homogenous spherical shell of mass with a radius $r_1 > 2m$, concentric with the central point mass m. The mass of the shell shall be m_1, its inner boundary r_1 and its thickness Δr_1. The presence of a massive shell at $r > r_1$ modifies the energy momentum tensor $T_{\mu\nu}$, introducing an additional step function at $r = r_1$. Yet, $T_{\mu\nu}$ still vanishes for $0 < r < r_1$.

We apply the new boundary condition by reevaluating Schwarzschild's original analysis. His differential equations for the line element

$$\mathrm{d}s^2 = -f_0 \mathrm{d}x_0^2 + f_1 \mathrm{d}x_1^2 + f_2 \frac{\mathrm{d}x_2^2}{1 - x_2^2} + f_3 \left(1 - x_2^2\right) \mathrm{d}x_3^2 \tag{7}$$

in his transformed coordinates

$$x_0 = ct\ , \quad x_1 = \frac{r^3}{3}\ , \quad x_2 = -\cos\vartheta\ , \quad x_3 = \varphi \tag{8}$$

remain valid for $0 < r < r_1$, where f_0, f_1, and $f_2 = f_3$ are arbitrary functions of x_1. For the examination of Equations 7, 8 we need the boundary condition at $x_{1,1} = r_1^3/3$,

$$f_2(x_{1,1}) = f_3(x_{1,1}) = r_1^2 = (3x_{1,1})^{2/3}, \tag{9}$$

which replaces the original boundary condition $f_2 = f_3 = r^2$ for $r \to \infty$. Schwarzschild formulated the field equations as.

$$\partial_1\left(\frac{\partial_1 f_1}{f_1}\right) = \frac{1}{2}\left(\frac{\partial_1 f_1}{f_1}\right)^2 + \left(\frac{\partial_1 f_2}{f_2}\right)^2 + \frac{1}{2}\left(\frac{\partial_1 f_0}{f_0}\right)^2 \tag{10}$$

$$\partial_1\left(\frac{\partial_1 f_2}{f_1}\right) = 2 + \frac{(\partial_1 f_2)^2}{f_1 f_2} \tag{11}$$

$$\partial_1\left(\frac{\partial_1 f_0}{f_1}\right) = \frac{(\partial_1 f_0)^2}{f_0 f_1} \tag{12}$$

$$f_0 f_1 f_2^2 = 1 \ \text{ or } \ \frac{\partial_1 f_0}{f_0} + \frac{\partial_1 f_1}{f_1} + 2\frac{\partial_1 f_2}{f_2} = 0\ . \tag{13}$$

Equation 12, representing the time and energy related gravitational component $\mu = \nu = 0$, would include a term $T_{00}(x_1)$ for $x_1 > x_{1,1}$. However, the additional non-vanishing values only contribute beyond the boundary $x_{1,1}$ and need not be included in the integration, similar to the mass point $T_{00}(x_1 = 0)$. The functions f_μ are continuous and $T_{\mu\nu}$ vanishes in the entire integration interval $0 < x_1 < x_{1,1}$.

By factoring out f_0/f_1 from the left-hand differential of Equation 12 transforms to

$$\partial_1\left(\frac{\partial_1 f_0}{f_0}\right) = \frac{\partial_1 f_0\, \partial_1 f_1}{f_0 f_1}. \tag{14}$$

Dividing Equation 14 by $\partial_1 f_0/f_0$ enables its integration to

$$\frac{\partial_1 f_0}{f_0} = \alpha f_1 \tag{15}$$

with the integration constant α, as in Schwarzschild's analysis. Adding Equation 10 and Equation 14 while considering Equation 13 results in

$$-2\,\partial_1\left(\frac{\partial_1 f_2}{f_2}\right) = 3\left(\frac{\partial_1 f_2}{f_2}\right)^2, \tag{16}$$

which can be integrated to

$$\frac{\partial_1 f_2}{f_2} = \frac{2}{3x_1 + \rho} \ \Rightarrow\ f_2 = \lambda(3x_1 + \rho)^{2/3} \tag{17}$$

with the integration constants λ and ρ. The boundary condition in Equation 9 indicates $\lambda = 1$. Unlike Schwarzschild's original boundary condition at infinity, it in addition requires $\rho = 0$, that is,

$$f_2 = (3x_1)^{2/3}. \tag{18}$$

From Equations 13, 15, 18 we find

$$\partial_1 f_0 = \alpha f_2^{-2} = \alpha(3x_1)^{-4/3}. \tag{19}$$

Integrating Equation 19 leads to

$$f_0 = -\alpha(3x_1)^{-1/3} + \mu \tag{20}$$

with the integration constant μ. In contrast to the original boundary condition at $r \to \infty$, μ cannot be determined by Equation 9. Equations 13, 18 and 20 combine to

$$f_1 = \frac{1}{f_0 f_2^2} = \frac{(3x_1)^{-4/3}}{\mu - \alpha(3x_1)^{-1/3}}. \tag{21}$$

In Schwarzschild's original analysis, Equation 11 was immediately satisfied. Here we use it to determine the integration constant μ:

$$6\mu - 4\alpha(3x_1)^{-1/3} = 2 + 4\mu - 4\alpha(3x_1)^{-1/3} \ \Rightarrow \mu = 1\ . \tag{22}$$

Equations 18, 20, 21 result in Schwarzschild's original solution

$$f_0 = 1 - \frac{\alpha}{r}\ , \quad f_1 = \frac{1}{r^4 f_0}\ , \quad f_2 = f_3 = r^2\ , \tag{23}$$

which for $\alpha = 2m$ in the usual polar coordinates leads to the Schwarzschild coordinates

$$\mathrm{d}s^2 = -\left(1 - \frac{2m}{r}\right)\mathrm{d}t^2 + \left(1 - \frac{2m}{r}\right)^{-1}\mathrm{d}r^2 + r^2\mathrm{d}\Omega^2, \quad r < r_1. \tag{24}$$

Equation 24 for the point mass m within the massive shell m_1 is identical to Schwarzschild's original solution, regardless of the surrounding shell. According to the Birkhoff theorem, any spherically symmetric mass distribution in vacuum may be represented by a point mass at $r = 0$. Thus, the energy momentum configuration including the shell m_1 may in the region $r \geq r_1 + \Delta r_1$ be regarded as a point mass at the origin. The same solution of Equation 24 results, however, for the total mass $m \to m + m_1$. The outer solution converges to the inner one for $\Delta r_1 \to 0$ ($m_1 \to 0$). An observer located at $2m < r < r_1$ experiences a gravitational field that does not change if the massive shell is removed. Only the inner mass m contributes to the gravitational field. The shell's own gravitational field thus vanishes at $r < r_1$.

This finding extends the validity of the relativistic Newtonian shell theorem to the full scope of its classical equivalent. Formerly, it was valid due to Birkhoff's theorem, yet, only for shells with pure vacuum inside. According to the findings above the gravitational field of a spherically symmetric shell vanishes inside, regardless of the radial mass distribution inside or across the shell.

5 Black hole formation

In the section on hand, we apply the results above to a contracting star experiencing a material collapse.

The Tolman-Oppenheimer-Volkoff equation for the pressure $p(r)$ within a spherically symmetric, perfect fluid in hydrostatic equilibrium describes the formation of black holes. Its only closed solution is the one for a constant energy density $\rho_e(r) = c^2\rho_0$, where ρ_0 is the mass density of the fluid,

$$p(r) = c^2\rho_0 \frac{f(r) - f(r_0)}{3f(r_0) - f(r)}, \quad f(r) = \sqrt{1 - \frac{r_{S0}r^2}{r_0^3}} \tag{25}$$

with the radius of the star r_0 and its Schwarzschild radius $r_{S0} = 2Gm_0/c^2$, where m_0 denotes the total mass of the star, G the gravitational constant and c the speed of light. The hydrostatic pressure $p(r)$ has its maximum in the center, $r = 0$, and turns singular for $3f(r_0) \to 1$. Thus, for hydrostatic equilibrium, r_0 must meet the constraint $f(r_0) > 1/3$, which corresponds to $r_0 > 9r_{S0}/8$. The diverging pressure leads to a collapse of the material structure when the star contracts to a radius $r_0 \approx 9r_{S0}/8$. The collapse starts in the center, within an infinitesimal radius $\mathrm{d}r = \varepsilon r_0$ ($\varepsilon \ll 1$). The collapsing mass may be expressed as

$$\mathrm{d}m = \rho_0 \frac{4\pi}{3}\mathrm{d}r^3 = \varepsilon^3 m_0. \tag{26}$$

The corresponding Schwarzschild radius $\mathrm{d}r_S \propto \mathrm{d}m$ of the collapsing region is

$$\mathrm{d}r_S = \varepsilon^3 r_{S0} \approx \frac{8\varepsilon^2}{9}\mathrm{d}r \ll \mathrm{d}r. \tag{27}$$

According to the hitherto notion, all the stellar fluid falls into the singularity. The findings summarized in Section 3.3 indicate that instead, the fluid of the adjacent infinitesimal spherical shell falls onto the event horizon at $\mathrm{d}r_S \ll \mathrm{d}r$ from Equation 27. The results from the preceding Section 4 allow for a black hole within a spherical shell of mass, that is, for a black hole within a black hole. Each of the subsequently falling shells $4\pi r^2\mathrm{d}r$ increases $r_S(r)$ such that

$$r_S(r) = r_{S0}\frac{r^3}{r_0^3} \approx \frac{8r^3}{9r_0^2}. \tag{28}$$

The monotonous increase in r of Equation 28 ensures $r_S(r) < r$ for $0 < r \leq r_0$. Finally, the star has immediately contracted from its previous radius r_0 to the new radius $r_{S0} \approx 8r_0/9$ and results in a mass density distribution

$$\rho(r) = \frac{1}{4\pi r^2}\frac{\mathrm{d}m(r_S)}{\mathrm{d}r_S} = \frac{c^2}{8\pi G r^2}. \tag{29}$$

The density distribution (Equation 29) increases with decreasing radius, as depicted in Figure 3. Each concentric shell $4\pi r^2\mathrm{d}r$ by the locally diverging curvature of space-time is caught within itself. Energy flow is freely possible in the angular directions but, as seen from a distant viewpoint, asymptotically retains in the radial direction. The density distribution $\rho(r)$ is independent of the star's mass m_0, but integrating it from the center to the star's new radius r_{S0} results in m_0, as expected. The mass density $\rho(r)$ diverges for $r \to 0$, which agrees with the collapse of the central region. A natural measure for the size of the collapsed region is the Planck density $\rho_P = c^5/(\hbar G^2)$. The radius of the collapsed sphere by Equation 29 may be estimated as $r_k = c/\sqrt{8\pi G\rho_P} = \ell_P/\sqrt{8\pi}$, about a factor five below the Planck length $\ell_P = \sqrt{\hbar G/c^3}$, the limit where the established physical concepts become invalid. We may conclude that all the stellar fluid is preserved.

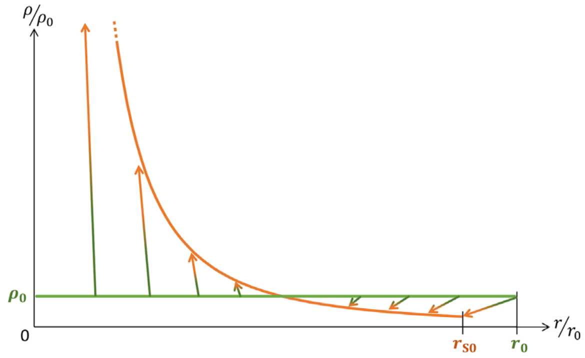

FIGURE 3
Graph of the density distribution $\rho(r)$ (orange line), diverging for $r \to 0$, in relation to the original constant density ρ_0 (green line). The green-orange arrows schematically illustrate how the stellar fluid drops from the state ρ_0 to $\rho(r)$.

These findings lead to a notion of black holes as stratified bodies of fluid with a well-defined density distribution and material structure, except for a collapsed central region with a size below the limits of the approved concepts of physics. During the accretion of additional mass, each new infinitesimal layer is caught in the locally diverging curvature of the space-time and for the next layer increases the Schwarzschild radius of the black hole.

6 Discussion and conclusion

The notion of Schwarzschild-Kerr black holes as engulfing abysses remained unchallenged for decades. In the past few years, however, several teams noted the weaknesses of this concept and proposed alternate models, trying to avoid or circumvent the singularity at the event horizon. Mimickers, gravastars, boson stars, or gravitational metamaterials are examples. The study on hand bases on the Schwarzschild metric alone, with no additional assumptions whatsoever.

We fathomed the conditions of the symmetry group of the spherically symmetric metric and identified an inherent discontinuity at the event horizon, the Birkhoff theorem being locally inapplicable. The spherically symmetric solutions form two subclasses, locally at the event horizon belonging to separate manifolds. The irregular Schwarzschild solution on the one hand represents a physical space-time. The globally regular solutions on the other, if they continuously extend the metric beyond the event horizon, violate the conditions of the symmetry group. As a consequence, the event horizon is impassable from the outside.

A revaluation of the Schwarzschild solution, integrated within a concentric shell of mass as a boundary, resulted in exactly the same coordinates as in vacuum to infinity. The Newtonian shell theorem was found fully valid in the context of the Schwarzschild metric. In combination with the impassable event horizon the new Schwarzschild solution allowed reconsidering the formation of black holes. We are thus facing a new notion of black holes as a stratification of stellar fluid, formed by a continuous progression of event horizons. Black holes in

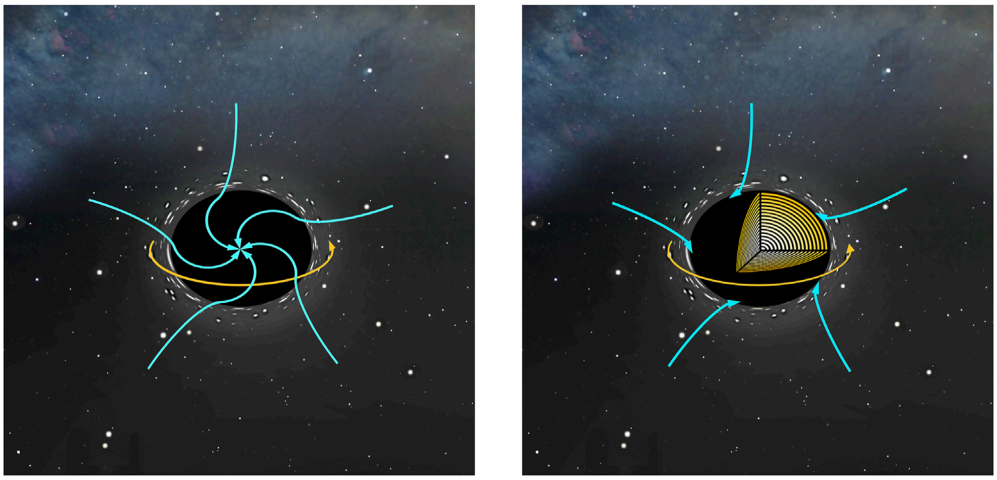

FIGURE 4
Artwork of matter falling into a rotating Kerr black hole (left, obsolete vision) and viewed as a black star according to the present analysis (right). Matter falling into the gravitational field is sketched in blue, the black star's fluid and the rotation in orange.

fact are black stars. The central singularity remains but the material structure is preserved all the way down to the Planck length, where the approved physical concepts become invalid.

Chou [15] presented a method to construct from the Schwarzschild coordinates the Kerr metric for rotating masses. He transforms the coordinates to the new symmetry and then adds the rotational energy. We conclude that equivalent restrictions apply to the Kerr metric, as to the Schwarzschild metric. Figure 4 depicts an artwork of a Kerr black hole according to the former and the current vision. The concept discussed above only limits the radial motion of stellar fluid, while angular motion is freely possible. Thus, the notion of black stars is applicable to rotating gravitational objects as well. The current vision is more plausible than the hitherto conception of Kerr black holes, because the angular momentum is naturally contained in the rotating mass and does not need to be attributed to a rotating vacuum.

The concept of black stars completely and unavoidably replaces the one of black holes. Yet, these findings affect the interior of black stars only. Outside the event horizon, all former studies remain valid as before. Time still dynamically interacts with space and the problem of time in quantum gravity as summarized by Anderson [16] remains. Yet, even upon gravitational singularities the timeline always remains time. This indicates a conceptual shift closer towards quantum theories, confirming, for example, the general suitability of the standard approach to canonical quantum gravity, that is, decomposing the space-time into a foliation of spacelike hypersurfaces as described, for instance, by Kiefer [17].

The material structure within a black star is preserved. As seen by a distant observer, each infinitesimal concentric layer is caught within itself, allowing for angular movements only. Each layer keeps on falling freely, even though it stalls with respect to the coordinate radius. The radial component of its proper line element diverges, while the angular components asymptotically approach the well-defined limit of the local event horizon. A falling observer thus experiences an ever-increasing volume. There is no limitation regarding the energy density as observed from a distant stationary viewpoint. The information paradox and the potential conflict with the Pauli exclusion principle no longer need be discussed. They simply become irrelevant. These results might, for instance, significantly affect black hole thermodynamics. The consequences may be far-reaching, even regarding cosmology and the evolution of the universe.

Data availability statement

The original contributions presented in the study are included in the article/supplementary material, further inquiries can be directed to the corresponding author.

Author contributions

JB-M: Formal Analysis, Visualization, Resources, Validation, Conceptualization, Methodology, Writing – review and editing, Investigation, Writing – original draft, Software.

Funding

The author(s) declare that no financial support was received for the research and/or publication of this article.

Conflict of interest

The author declares that the research was conducted in the absence of any commercial or financial relationships that could be construed as a potential conflict of interest.

Generative AI statement

The author(s) declare that no Generative AI was used in the creation of this manuscript.

Publisher's note

References

1. Schwarzschild K. On the gravitational field of a mass point according to Einstein's theory. *Sitzungsber Preuss Akad Wiss* (1916) 189–96. doi:10.48550/ARXIV.PHYSICS/9905030

2. Painlevé P. La mécanique classique et la théorie de la relativité. *C R Acad Sci (Paris)* (1921) 173:677–80.

3. Jebsen JT. Über die allgemeinen kugelsymmetrischen Lösungen der Einsteinschen Gravitationsgleichungen. *Ark Matem., Astron Fys* (1921) 15:1–9.

4. Kruskal MD. Maximal extension of the Schwarzschild metric. *Phys Rev* (1960) 119:1743–5. doi:10.1103/PhysRev.119.1743

5. Szekeres G. On the singularities of a Riemannian manifold. *Publ Math Debrecen* (1960) 7:285–301. doi:10.5486/pmd.1960.7.1-4.26

6. Lemos JPS, Zaslavskii OB. Black hole mimickers: regular versus singular behavior. *Phys Rev D* (2008) 78:024040. doi:10.1103/PhysRevD.78.024040

7. Casadio R, Kamenshchik A, Ovalle J. From black hole mimickers to black holes. *Phys Rev D* (2024) 109:024042. doi:10.1103/PhysRevD.109.024042

8. Mazur PO, Mottola E. Gravitational condensate stars: an alternative to black holes. *Universe* (2023) 9:88. doi:10.3390/universe9020088

9. Tkachev II. Coherent scalar-field oscillations forming compact astrophysical object. *Sov Astron Lett* (1986) 305.

10. Luongo O. Gravitational metamaterials from optical properties of spacetime media (2025). Available online at: https://arxiv.org/abs/2504.09987.

11. Yang F, Zhang Z, Xu L, Liu Z, Jin P, Zhuang P, et al. Controlling mass and energy diffusion with metamaterials. *Rev Mod Phys* (2024) 96:015002. doi:10.1103/RevModPhys.96.015002

12. Plebański J, Krasiński A. *An introduction to general relativity and cosmology*. New York: Cambridge Univ. Press (2006). doi:10.1017/CBO9780511617676

13. Kerr RP. Gravitational field of a spinning mass as an example of algebraically special metrics. *Phys Rev Lett* (1963) 11:237–8. doi:10.1103/PhysRevLett.11.237

14. Schwarzschild K. On the gravitational field of a sphere of incompressible fluid according to Einstein's theory. *Sitzungsber Preuss Akad Wiss* (1916) 424–34. doi:10.48550/ARXIV.PHYSICS/9912033

15. Chou Y. A radiating Kerr black hole and Hawking radiation. *Heliyon* (2020) 6:e03336. doi:10.1016/j.heliyon.2020.e03336

16. Anderson E. Problem of time in quantum gravity. *Ann Phys* (2012) 524:757–86. doi:10.1002/andp.201200147

17. Kiefer C. Concept of time in canonical quantum gravity and string theory. *J Phys Conf Ser* (2009) 174:012021. doi:10.1088/1742-6596/174/1/012021

18. Liebling SL, Palenzuela C. Dynamical boson stars. *Living Rev Relativ* (2017) 26:1. doi:10.1007/s41114-023-00043-4