Lorentzian wormholes in Einstein-Gauss-Bonnet theory

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Lorentzian wormhole solutions are investigated in the context of the D-dimensional Einstein-Gauss-Bonnet theory of gravitation. These wormholes are found to have features depending on the dimensionality of the spacetime and the coupling coefficient α of the Gauss-Bonnet combination. In a large number of cases, the wormhole throat radius is constrained to have a value greater than a certain number depending on D and α. The possibility of obtaining solutions with normal and exotic matter limited to the vicinity of the throat is explored. Similar to the situation in general relativity, the violation of the weak energy condition persists for $\alpha > 0$. For $\alpha < 0$, this condition may or may not be violated depending on the nature of an inequality involving $|\alpha|$, D, the radius r, and the wormhole shape function $b(r)$.

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I. INTRODUCTION

In the last few years, wormholes have rapidly grown into an active area of research. Initially, the stress was primarily on Euclidean wormholes which modeled topology changing processes in quantum gravity by means of quantum tunneling [1]. However, with the pioneering work of Morris and Thorne [2] and Morris, Thorne, and Yurtsever [3], it became clear that there exist certain solutions of the Einstein field equations of general relativity which possess wormholelike features. Unfortunately, it was found that the matter that threaded the wormhole violated the weak energy condition (WEC) near the throat. Morris and Thorne [2] argued that such a violation of the WEC was permissible in the context of quantum field theory in curved spacetime and its theoretical procedure for discarding infinities of the zero-point energy from the stress tensor. Based on their suggestion [2], Hochberg and Kephart [4] have recently obtained Lorentzian wormholes using matter constructed out of squeezed vacuum states of light having stress tensors which, when coupled to gravity, can violate the WEC. On the other hand, Visser [5], while discussing the construction of such Lorentzian wormholes using "Schwarzschild surgery" and the junction-condition formalism, argued that the WEC cannot be accepted as a worthwhile principle because experiments have ruled out its validity.

There also have been quite a few papers discussing Lorentzian wormholes in alternative theories of gravity. Hochberg [6] has discussed such solutions in the context of $\mathcal{R}+\mathcal{R}^2$ theories in four dimensions. Moffat [7] has recently shown that the violation of the WEC persists in his nonsymmetric theory of gravitation.

In this paper we discuss Lorentzian wormholes in the Einstein-Gauss-Bonnet (EGB) theory of gravity. The action for this theory consists of the usual Einstein-Hilbert term plus the Gauss-Bonnet (GB) combination. In four dimensions the EGB theory reduces to general relativity. This is due to the fact that the GB combination reduces to a pure divergence in four dimensions by virtue of the Bach identity. Spherically symmetric black-hole and cosmological solutions of the field equations of the EGB theory have been discussed in detail by various authors [8,9]. Higher-dimensional Euclidean wormholes in this theory have been studied by Gonsalez-Diaz [10], and Jianjun and Sicong [11]. Our focus here is on Lorentzian wormholes. We shall use dimensionally extended versions of the metric ansatz and the stress-energy tensor used by Morris and Thorne [2]. The presence of extra dimensions as well as the GB combination leads to the existence of a wide class of solutions. The WEC is violated here also but for the case in which the coupling coefficient for the GB combination is negative and satisfies a certain inequality involving the wormhole shape function. In the end we shall suggest a construction of a solution with matter satisfying the WEC everywhere.

The paper is organized as follows. In Sec. II we discuss the field equations of the EGB theory. Section III deals with the various kinds of solutions together with the construction mentioned at the end of the previous paragraph. Section IV is a summary of the results.

II. THE FIELD EQUATIONS

The action integral for the EGB is given as

$$I=\int d^Dx\sqrt{-g}\,[\kappa\mathcal{R}+\alpha(\mathcal{R}_{\mu\nu\rho\sigma}\mathcal{R}^{\mu\nu\rho\sigma}-4\mathcal{R}_{\mu\nu}\mathcal{R}^{\mu\nu}+\mathcal{R}^2)]+I_{\text{matter}}\ , \tag{1}$$

where D is the dimensionality of spacetime and $\kappa>0$. Henceforth, we choose to work in units where $c=\kappa=1$. The field equations that follow from such an action are given as

$$\begin{aligned}0=&[\mathcal{R}_{\mu\nu}-\tfrac{1}{2}g_{\mu\nu}\mathcal{R}]\\&-\alpha[\tfrac{1}{2}g_{\mu\nu}(\mathcal{R}_{\alpha\beta\gamma\delta}-\mathcal{R}^{\alpha\beta\gamma\delta}-4\mathcal{R}_{\alpha\beta}\mathcal{R}^{\alpha\beta}+\mathcal{R}^2)\\&-2\mathcal{R}\mathcal{R}_{\mu\nu}+4\mathcal{R}_{\mu\alpha}\mathcal{R}^{\alpha}_{\nu}+\mathcal{R}_{\alpha\beta}\mathcal{R}^{\alpha\beta}_{\mu\nu}\\&-2\mathcal{R}_{\mu\alpha\beta\gamma}\mathcal{R}^{\alpha\beta\gamma}_{\nu}]-T_{\mu\nu}\ .\end{aligned} \tag{2}$$

The α which appears above as the coupling coefficient of the GB combination is positive as long as we consider the EGB theory as the low-frequency limit of superstring theory [12,13]. It is important to note that if the EGB theory is assumed as a theory on its own right (it is in fact the first-order correction to general relativity suggested by Lovelock [13,14]) there is no restriction on the sign of α. We shall take α with values in various ranges and derive the consequences. A similar stand regarding α has been taken by Wiltshire [8] and Wheeler [8] while discussing the spherically symmetric black hole and cosmological solutions of the EGB theory.

Our metric ansatz is the same as Morris and Thorne, except that the two-sphere is replaced by a $(D-2)$-sphere. It is given as

$$ds^2=-e^{2\phi(r)}dt^2+\left[1-\frac{b(r)}{r}\right]^{-1}dr^2+r^2d\Omega_{D-2}^2 . \tag{3}$$

Here $\phi(r)$ and $b(r)$ are the redshift function and the shape function, respectively. $d\Omega_{D-2}^2$ is the metric on the surface of a $(D-2)$-sphere.

We can write Eq. (3) in the proper orthonormal basis as

$$ds^2=-e^0\otimes e^0+e^1\otimes e^1+\sum_{i=2}^{D-1}e^i\otimes e^i , \tag{4}$$

where the basis one-forms are

$$\begin{aligned} &e^0=e^{\phi}dt , \\ &e^1=[1-b/r]^{-1/2}dr , \\ &e^2=r\,d\theta_2 , \\ &\vdots \\ &e^i=r\sin\theta_2\ldots\sin\theta_{i-1}d\theta_i , \\ &\vdots \\ &e^{D-1}=r\sin\theta_2\ldots\sin\theta_{D-2}d\theta_{D-1} . \end{aligned} \tag{5}$$

In our notation, there is no θ_1. The angular coordinates are denoted by θ_i, where $i=2,3,\ldots,D-1$.

Using the Cartan equations of structure we can derive the curvature two-forms:

$$\begin{aligned} &\mathcal{R}^0_1=Me^0\wedge e^1 , \quad \mathcal{R}^1_i=Pe^1\wedge e^i , \\ &\mathcal{R}^0_i=Ne^0\wedge e^i , \quad \mathcal{R}^i_j=Qe^i\wedge e^j . \end{aligned} \tag{6}$$

For any fixed i, j can have values $i<j\leq(D-1)$. M,N,P,Q are

$$M=\left[1-\frac{b}{r}\right]\left[-\phi''-\phi'^2+\frac{\phi'(b'r-b)}{2r(r-b)}\right] , \tag{7a}$$

$$N=-\frac{\phi'}{r}(1-b/r) , \tag{7b}$$

$$P=\frac{1}{2r^3}(b'r-b) , \tag{7c}$$

$$Q=\frac{b}{r^3} , \tag{7d}$$

where the prime denotes differentiation with respect to r. From these curvature two-forms all components of the full Riemann-Christoffel curvature tensor and hence the Ricci tensor and the Ricci scalar can be derived.

The stress-energy tensor in the static observer's frame is given as

$$T_{00}=\rho(r) , \quad T_{11}=-\tau(r) , \quad T_{ij}=p(r)\delta_{ij} , \tag{8}$$

where $i,j=2,3,\ldots,D-1$.

The field equations turn out to be the following:

$$\begin{aligned} \rho(r)=(D-2)&\left[\left[\frac{D-3}{2}\right]Q+P\right] \\ &+\bar{\alpha}Q(D-2)\left[\left[\frac{D-5}{2}\right]Q+2P\right] , \end{aligned} \tag{9}$$

$$\begin{aligned} \tau(r)=(D-2)&\left[\left[\frac{D-3}{2}\right]Q+N\right] \\ &+\bar{\alpha}Q(D-2)\left[\left[\frac{D-5}{2}\right]Q+2N\right] , \end{aligned} \tag{10}$$

$$\begin{aligned} p(r)=-M-(D-3)N-(D-3)P-\tfrac{1}{2}(D-3)(D-4)Q& \\ -\tfrac{1}{2}(D-5)(D-6)\bar{\alpha}Q^2& \\ -4(D-5)\bar{\alpha}NP-2(D-5)\bar{\alpha}PQ& \\ -2(D-5)\bar{\alpha}NQ-2\bar{\alpha}MQ ,& \end{aligned} \tag{11}$$

where $\bar{\alpha}=(D-3)(D-4)\alpha$. One can check that for $D=4$, Eqs. (9)–(11) reduce to the Einstein field equations as derived in the paper by Morris and Thorne [2]. For $\alpha=0$ and $D>4$ the field equations are those for higher-dimensional general relativity. Also, the vacuum solutions of the field equations (9)–(11) give the standard Boulware-Deser black-hole spacetime [9].

III. THE SOLUTIONS

A. General constraints on Lorentzian wormholes

Morris and Thorne [2] have discussed in detail the general constraints that need to be obeyed if the spacetime is to have wormhole-like features. We mention them here for the sake of completeness.

Constraint 1: At the throat of the wormhole, $r=b(r)=b_0$, b_0 being the minimum value of r.

Constraint 2: $b(r)/r\leq 1$ throughout the spacetime. This is required to ensure the finiteness of the proper radial distance defined by

$$l(r)=\pm\int_{b_0}^{r}\frac{dr}{[1-b/r]^{1/2}} . \tag{12}$$

The $\pm$ signs refer to the two asymptotically flat regions which are connected by the wormhole. The equality sign in $b(r)/r \leq 1$ holds only at the throat.

Constraint 3: As $l \to \pm\infty$ (or, $r \to \infty$), $b(r)/r \to 0$. This is the asymptotic flatness condition on the wormhole spacetime.

Constraint 4: $\phi(r)$ is finite throughout the spacetime to ensure the absence of horizons and singularities.

In this paper we are not interested in discussing the traversibility constraints mentioned by Morris and Thorne [2].

B. Exoticity of matter near the throat

In four-dimensional general relativity it was shown in Ref. [2] that the flaring-out condition near the throat of the wormhole led to the fact that $\tau_0 > \rho_0$, where τ_0 and ρ_0 are the values of τ and ρ near the throat. The WEC ($T_{\mu\nu}u^\mu u^\nu \geq 0$, where u^μ is the timelike four-velocity of the observer), which requires $\tau < \rho$, is therefore violated near the throat. Such matter was termed as "exotic." From the field equations for the EGB theory we get the following expression for $(\tau - \rho)$:

$$\tau - \rho = (D-2)[1+2\bar{\alpha}Q][N-P] \ . \tag{13}$$

The ξ function defined by Morris and Thorne [2] is $(\tau-\rho)/|\rho|$. The essential difference between the expression for $\tau - \rho$ in four-dimensional general relativity (GR) and the one for the EGB theory is the presence of two extra factors, $(D-2)$ and $[1+2\bar{\alpha}Q]$. The former factor originates from the dimensionality of the spacetime; the latter one is present only in the EGB theory. The factor $[N-P]$ is exactly identical to the expression for $\tau-\rho$ in GR. We examine $\tau-\rho$ near the throat. This leads to

$$\tau_0 - \rho_0 = (D-2)\left[1+\frac{2\bar{\alpha}}{b_0^2}\right][N-P]_{r=b=b_0} \ . \tag{14}$$

Now $[N-P]_{r=b=b_0} > 0$ as shown by Morris and Thorne [2]. Thus, for an EGB-theory wormhole, matter near the throat is exotic or normal if the quantity $(1+2\bar{\alpha}/b_0^2)$ is positive or negative, respectively. The above-stated conditions can be thought of as constraints on b_0 or α. We prefer to choose α and let this choice determine b_0. If $\alpha > 0$, the quantity $(1+2\bar{\alpha}/b_0^2)$ is always positive and the WEC is violated near the throat. Also, the mere fact that $(1+2\alpha/b_0^2) > 0$ does not imply any lower bound on b_0. For α negative, the quantity $(1-2|\bar{\alpha}|/b_0^2)$ can be positive or negative. If it is positive, then $b_0 > (2|\bar{\alpha}|)^{1/2}$. Otherwise, $b_0 < (2|\bar{\alpha}|)^{1/2}$. Thus, matter can be normal near the throat but, as will be shown below, other constraints forbid the existence of a solution with normal matter everywhere.

The next obvious question to ask is whether it is possible to obtain solutions in the EGB theory for which matter is exotic or normal everywhere. One can also investigate the possibility of having solutions with exotic or normal matter confined to the region in the vicinity of the throat. In the following subsections these situations are discussed.

C. Solutions with $\alpha > 0$ and $\tau > \rho > 0$, $\phi = 0$

From expression (13) we noticed that if $\alpha > 0$, $\phi = 0$, then $\tau > 0$ everywhere. This implies that $b' < b/r$ everywhere. Equation (9), after some rearrangements, reduces to

$$\rho(r) = (D-2)\frac{X}{2r^3} + (D-2)\bar{\alpha}b(r)\frac{Y}{2r^6} \ , \tag{15}$$

where

$$X = (D-4)b + b'r \ ,$$
$$Y = (D-7)b + 2b'r \ . \tag{16}$$

Thus $\rho(r)$ can be greater than zero if any of the following hold:

(a) $X = 0$, $Y > 0$,

(b) $X < 0$, $Y > 0$; $|X| < \bar{\alpha}bY/r^3$,

(c) $X > 0$, $Y = 0$,

(d) $X > 0$, $Y > 0$,

(e) $X > 0$, $Y < 0$; $X > \bar{\alpha}b|Y|/r^3$.

Thus, an acceptable zero-tidal-force wormhole solution in this case has to obey the condition $b' < b/r$, the general constraints described in Sec. III A and any one of the above conditions which lead to $\rho > 0$. For the cases (a) and (b) a little amount of analysis will show that we require $D < 1$ for $\rho > 0$, which is impossible. We shall study the remaining three cases below. Our strategy is to begin with the validity of the $\rho > 0$ condition. Then we examine the proposed solutions for the other constraints.

Case (c): In this case $Y = 0$ leads to the differential equation

$$b' = -\frac{(D-7)}{2}\frac{b}{r} \ .$$

The unique wormhole solution is

$$b(r) = b_0^{(D-5)/2} r^{-(D-7)/2} \ . \tag{17}$$

The solution in $D = 5$ is ruled out since it violates the general constraint (2).

Case (d): The conditions on X and Y give the following inequality:

$$b' > -\frac{(D-7)}{2}\frac{b}{r} \ .$$

In five dimensions this becomes $b' > b/r$, whereas $\tau > \rho$ led to $b' < b/r$. Thus, in five dimensions this kind of solution is ruled out. For $D \geq 6$ a large variety of solutions are possible. Some examples are the following:

Power-law solution

$$b(r) = b_0^{m/D} r^{1-m/D} \ ; \quad 0 < m < \frac{D(D-5)}{2} \ . \tag{18}$$

Logarithmic solution

$$b(r) = \frac{r}{\ln r}\ln b_0 \ ; \quad b_0 > e^{2/(D-5)} \ . \tag{19}$$

Hyperbolic solution

$$b(r)=\frac{b_0}{\tanh b_0}\tanh r \ . \tag{20}$$

However, this solution is valid only for $D\geq 7$. For $D=6$ it is valid only if $\sinh 2r<4r$. This implies limiting the range of r up to a certain value r_0. Consequently, no solution is possible with exotic matter everywhere.

Case (e): The conditions $X>0$ and $Y<0$ imply

$$-\left[\frac{D-7}{2}\right]>\frac{b'r}{b}>-(D-4) \ .$$

An example of a solution obeying all conditions is

$$b(r)=b_0^{m/D}r^{1-m/D} \ , \tag{21}$$

where m is always positive and takes values such that

$$\frac{D(D-5)}{2}<m<D(D-3) \ . \tag{22}$$

However, the inequality between X and Y leads to

$$b_0>\left[12\alpha\frac{m}{2D-m}\right]^{1/2} , \tag{23}$$

which requires $m<2D$. This new constraint on m along with Eq. (22) restricts D such that $4<D<9$. The throat radius b_0 is constrained to have a minimum value dependent on α, m, and D. The presence of the GB combination as well as extra dimensions is visible quite clearly here.

D. Solutions with $\tau>\rho>0$, $\phi=0$, and $\alpha<0$, $(1+2Q\bar{\alpha})>0$

Here we look for wormhole solutions with exotic matter everywhere but with $\alpha<0$. The condition $\tau>\rho$ near the throat implies

$$b_0>(2|\bar{\alpha}|)^{1/2} \ . \tag{24}$$

For general values of r, $b(r)$ should be such that $b(r)<r^3/(2|\bar{\alpha}|)$.

Just as in the previous section, the solutions in the present case should also satisfy the condition $b'<b/r$ and all the four general constraints. Also, to ensure the positivity of ρ, any one of the following five sets of conditions on X and Y should be satisfied:

(a) $X=0$, $Y<0$,

(b) $X<0$, $Y<0$; $|X|<b|\bar{\alpha}||Y|/r^3$,

(c) $X>0$, $Y=0$,

(d) $X>0$, $Y>0$; $X>b|\bar{\alpha}|Y/r^3$,

(e) $X>0$, $Y<0$.

Case (a): In this case, it can be shown that the unique wormhole solution has the form

$$b(r)=b_0^{D-3}r^{-(D-4)} \ . \tag{25}$$

Moreover, as stated above, the throat radius b_0 has a lower limit given by Eq. (24).

Case (b): The extra inequality relation between X and Y along with the conditions $b'r/b<1$ and $b/r<1$ leads to the fact that the domain of r is not $[b_0,\infty)$ but $[b_0,r_0)$. Thus, there exists in all cases an upper bound for r. This implies the nonexistence of solutions with exotic matter everywhere.

Case (c): Except for the existence of a lower bound for b_0 [Eq. (24)], this case is similar to case (c) of Sec. III C. The solution is given by Eq. (17) and is valid for $D\geq 6$.

Case (d): This case has a similarity with case (d) of Sec. III C. No solution is possible for $D=5$. All solutions mentioned there for $D\geq 6$ are applicable here. However, there will always be a lower bound on b_0. For every solution we have to choose the appropriate lower bound by comparing Eq. (24) and the extra inequality relation between X and Y. For example, in the power-law solution [Eq. (18)], we have to choose between Eq. (24) and the inequality:

$$b_0>\left[|\bar{\alpha}|\frac{D-5-2m/D}{D-3-m/D}\right]^{1/2} \ . \tag{26}$$

Since the right-hand side of Eq. (24) is greater than that of Eq. (26), the former one will give the lower limit for b_0.

Case (e): The only difference between this case and (e) of Sec. III is that the lower limit on b_0 here is determined by Eq. (24). Due to this, the power-law solution does not have the extra restriction $m<2D$. It is therefore valid for all $D>4$.

E. Solutions with $\rho>\tau$ and $\rho>0$

It is not possible to have wormhole solutions of the EGB field equations with normal matter everywhere. To have $\rho>\tau$, we need

$$b(r)>\frac{r^3}{2|\bar{\alpha}|} \ . \tag{27}$$

This, together with the fact that $b/r<1$, implies that any solution obtained in this case will be valid only up to a certain value of r, i.e., $r<(2|\bar{\alpha}|)^{1/2}$. A way out of this may be the following. Consider a solution with normal matter extended from the throat radius up to a certain radius r_c so that $b_0<r_c<(2|\bar{\alpha}|)^{1/2}$. At $r=r_c$ join the solution to the vacuum Boulware-Deser spacetime across a surface layer. However, this requires an extension of the junction-condition formalism of the GR to EGB theory.

F. Solutions with exotic matter limited to the throat

We can also have solutions with $\tau>\rho$, $\alpha>0$ such that the exotic matter is limited to the throat region. Following Ref. [2], we assume $r_a=b_0+\Delta r$, where Δr is the region of extension of exotic matter near the throat. To get a significant flaring-out from the throat, we need to have dz/dr at $r=r_a$ very near to 1, where $z(r)$ describes the $t=$ constant, $\theta=\pi/2$ section of the wormhole spacetime embedded in R^3. These lead to $b'<0$ once $\Delta r\ll b_0$. It is not possible to have $\rho>0$, $b'<0$, and $\Delta r\ll b_0$ simultaneously in four-dimensional GR. Let us consider the corresponding situation in the EGB theory for the case $X > 0$, $Y > 0$, and $\alpha > 0$. We require

$$\frac{|b'|r}{b} < \frac{1}{2}(D-7) \tag{28}$$

to have $\rho > 0$, with $b' < 0$. Combining this with

$$\frac{\Delta r}{b_0}(1+|b'|)=1 , \tag{29}$$

we get

$$\Delta r > \frac{2}{D-5} b_0 . \tag{30}$$

Thus, for $D > 7$, $\Delta r / b_0 < 1$ is not a contradiction. For higher dimensions, we see that Δr becomes increasingly smaller than b_0. The fact mentioned above is possible not only in the EGB theory but also in D-dimensional GR, where we have

$$\Delta r > \frac{1}{D-4} b_0 . \tag{31}$$

IV. CONCLUSION

In conclusion, it is important to figure out the distinguishing features of wormholes in the EGB theory as compared to the ones in GR.

(i) For $\tau > \rho > 0$ and $\alpha > 0$, an EGB wormhole exists in five dimensions only in the case in which the throat radius is constrained to have a minimum value dependent on α. Also there are solutions which have a minimum throat radius depending on D only (the logarithmic solution). Solutions with b_0 independent of D or α are also there; b_0 here can be arbitrary but finite. On the other hand, for $\tau > \rho > 0$ and $\alpha < 0$ with $1+2\bar{\alpha}Q > 0$, the solutions are forced to have a throat radius which is bounded below by a quantity dependent on D and α. Finally, if $\tau < \rho$, $\rho > 0$, $\alpha < 0$, and $1+2\bar{\alpha}Q < 0$, all solutions are defined only up to a certain value of $r = r_c$. Beyond $r = r_c$, normal matter cannot exist and the only way out, as suggested in this paper, is the type of solutions with normal matter for $b_0 < r < r_c - \epsilon$ and vacuum for $r > r_c - \epsilon$.

(ii) It has been shown in the EGB theory as well as in D-dimensional GR ($D > 5$) that one can limit exotic matter with $\rho > 0$ to an arbitrarily small region. This was not possible in four-dimensional GR.

(iii) The status of the WEC in the EGB theory, however, remains almost the same as it was in GR. For $\alpha > 0$, it is violated; for $\alpha < 0$ it may or may not be violated depending on whether $(1+2\bar{\alpha}Q)$ is greater or less than zero. Even if the WEC is not violated, one cannot construct EGB wormholes with normal matter everywhere.

It would be interesting to explore the possibility of constructing an EGB wormhole with the WEC satisfied everywhere. This implies investigating the solution with normal matter confined to the throat and vacuum elsewhere. We do not know whether the matching of the two different types of solutions would be possible at all.

Furthermore, one can investigate the existence and features of Lorentzian wormholes in other theories of gravitation. Einstein-Cartan-Sciama-Kibble theory and Brans-Dicke models are probable candidates for such analyses. String theory leads to other actions in lower dimensions with dilatons and rank-three antisymmetric tensor fields coupled to the usual Einstein-Hilbert term. Are there Lorentzian wormholes in such effective theories? Attempts at answering such questions and other related issues will be communicated in the future.

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