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A microwave realization of the Gaussian symplectic ensemble

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Following an idea by Joyner et al. [EPL, 107 (2014) 50004] a microwave graph with an antiunitary symmetry $T$ obeying $T^2 = -1$ has been realized. The Kramers doublets expected for such systems have been clearly identified and could be lifted by a perturbation which breaks the antiunitary symmetry. The observed spectral level spacings distribution of the Kramers doublets is in agreement with the predictions from the Gaussian symplectic ensemble, expected for chaotic systems with such a symmetry.

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Random matrix theory has proven to be an extremely powerful tool to describe the spectra of chaotic systems [1–4]. For systems with time-reversal symmetry (TRS) and no half-integer spin in particular there is an abundant number of studies, theoretical, numerical and experimental, showing that the universal spectral properties are perfectly well reproduced by the corresponding properties of the Gaussian orthogonal random matrix ensemble (GOE) (see Ref. 5 for a review). This is the essence of the famous conjecture by Bohigas, Giannoni and Schmit [3], which has received strong theoretical support, see e.g. Refs. 6–8. For systems with TRS and half-integer spin and systems with no TRS the Gaussian symplectic ensemble (GSE) and the Gaussian unitary ensemble (GUE), respectively, hold instead. There are three studies of the spectra of systems with broken TRS showing GUE statistics [9–11], all of them applying microwave techniques. For the GSE there is no experimental realization at all up to now. Only by using that a GSE statistics if the system is chaotic [20]. In addition it leads to Kramers degeneracy, i.e. application of $T$ to an energy eigenfunction yields an orthogonal eigenfunction with equal energy. A system with such a symmetry was recently found in the form of a quantum graph [21].

Quantum graphs were introduced by Kottos and Smilansky [22] to study various aspects of quantum chaos. The wave function on a quantum graph satisfies a one-dimensional Schrödinger equation on each of the bonds with suitable matching conditions (implying current conservation) at the vertices. Just, as for quantum billiards, there is a one-to-one mapping onto the corresponding microwave graph. This analogy has been used in a number of experiments including one on graphs with and without broken TRS [11, 23, 24].

To realize graphs with GSE symmetry the graph shown in Fig. 1(top) was proposed in [21]. It contains two geometrically identical subgraphs, but with phase shifts by $+\pi/2$ and $-\pi/2$, respectively, along two corresponding bonds. The two subgraphs are connected by a pair of bonds yielding a graph with a geometric inversion center. In addition there is another phase shift of $\pi$ along one of the two bonds, but not the other one! This is the crucial point: Due to this trick the total graph is symmetric with respect to an antiunitary operator $T$, squaring to minus one, $T^2 = -1$, where

$$T\psi(x_1) = +\psi^*(x_2), \quad T\psi(x_2) = -\psi^*(x_1),$$

i.e. if $\psi$ satisfies the Schrödinger equation as well as the vertex conditions then the same applies to $T\psi$. Here $x_1$ is a coordinate in subgraph 1, and $x_2$ the corresponding coordinate (related by inversion) in subgraph 2. Applying (1) twice shows $T^2 = -1$.

A complementary approach shall be given establishing a direct correspondence between the experiment and a spin 1/2 system. The wave function in a quantum graph
is subject to two constraints, continuity at the vertices and current conservation. In a microwave graph these constraints correspond to the well-known Kirchhoff rules governing electric circuits. They lead to a secular equation system having a solution only if the determinant of the corresponding matrix vanishes,

$$\text{det}[h(k)] = 0$$  \hspace{1cm} (2)

where the matrix elements of $h(k)$ are given by

$$h_{ij}(k) = \begin{cases} 
- \sum_{n \neq i} C_{in} \cot(\lambda_i \lambda_n) & i = j \\
C_{ij} e^{-\tau \varphi_{ij}} [\sin(\lambda_{ij})]^{-1} & i \neq j 
\end{cases}$$  \hspace{1cm} (3)

where $C_{ij} = 1$, if nodes $i$ and $j$ are connected, and $C_{ij} = 0$ otherwise. $\lambda_{ij}$ is the length of the bond connecting nodes $i$ and $j$. $\varphi_{ij}$ is a phase resulting, e.g., from a vector potential and breaks TRS if present. The equation holds for Neumann boundary conditions at all nodes, the situation found in the experiment. Details can be found in [25]. The solutions of the determinant condition (2) generate the spectrum of the graph.

Applied to the graph of Fig. 1 the secular matrix $h(k)$ may be written as

$$h = h_{\text{dis}} + v$$  \hspace{1cm} (4)

where $h_{\text{dis}}$ is the secular matrix for the disconnected subgraphs, and $v$ describes the connecting bonds. It is convenient to introduce an order of rows and columns according to $\{1, 2, \ldots, n; 1, 2, \ldots, \bar{n}\}$, where the numbers without bar refer to the vertices of subgraph 1, and the numbers with bar to those of subgraph 2. $h_{\text{dis}}$ may then be written as

$$h_{\text{dis}} = \begin{pmatrix} h_0 & \cdot \\ \cdot & h_0^* \end{pmatrix}$$  \hspace{1cm} (5)

where $h_0$ and $h_0^*$ are the secular matrices for each of the two subgraphs, respectively. Since the only difference between the subgraphs is the sign of the $\pi/2$ phase shift in one of the bonds, their secular matrices are just complex conjugates of each other, see Eq. (3). Assuming for the sake of simplicity that there is just one pair of bonds connecting node 1 with node 2, and node 1 with node 2, the matrix elements of $v$ are given by

$$v_{11} = v_{22} = v_{1\bar{1}} = v_{2\bar{2}} = -\cot(\lambda l)$$  \hspace{1cm} (6)

$$v_{12} = v_{21} = -v_{1\bar{2}} = -v_{2\bar{1}} = [\sin(\lambda l)]^{-1}$$  \hspace{1cm} (7)

$$v_{ij} = v_{\bar{i}j} = v_{ij} = v_{\bar{i}j} = 0 \quad \text{otherwise}$$  \hspace{1cm} (8)

where $l$ is the length of the bond connecting 1 with 2 and $\bar{1}$ with 2. The generalization to a larger number of bond pairs is straightforward.

Changing now the sequence of rows and columns to $\{1, \bar{1}; 2, \bar{2}; \ldots; n, \bar{n}\}$, the resulting $2n \times 2n$ matrix $h(k)$
may be written in terms of a \( n \times n \) matrix with quaternion matrix elements,

\[
[\tilde{h}(k)]_{nm} = [\text{Re}(h_0)]_{nm} + v_{nm} 1 - \text{Im}(h_0)_{nm} \tau_{z} - v_{nm} \tau_{y}
\]

where

\[
1 = \begin{pmatrix} 1 & \cdot \\ \cdot & 1 \end{pmatrix}, \quad \tau_{z} = \begin{pmatrix} -i & \cdot \\ \cdot & i \end{pmatrix}, \quad \tau_{y} = \begin{pmatrix} \cdot & -1 \\ 1 & \cdot \end{pmatrix}
\]

The determinant is not changed by this rearrangement of rows and columns,

\[
\det[\tilde{h}(k)] = \det[h(k)].
\]

The matrix elements \([\tilde{h}(k)]_{nm}\) commute with \(C\tau_{y}\), where \(C\) denotes the complex conjugate, and hence the whole matrix commutes with \(T = \text{diag}(C\tau_{y}, \ldots, C\tau_{y})\), where \(T\) squares to minus one, \(T^2 = -1\). This is exactly the situation found for spin 1/2 systems, and just as in such systems a two-fold Kramers degenerate spectrum is expected showing the signatures of the GSE provided the system is chaotic, see e.g. Chapter 2 of Ref. 4.

The requirements defined by Joyner et al. [21] to realize graphs with GSE symmetry pose some challenges. Since we did not know of a simple way to achieve phase shifts of \(\pm \pi/2\) along the bonds, we instead built two geometrically identical subgraphs, but with two circulators of opposite sense of rotation within the two subgraphs. A circulator is a T-shaped microwave device introducing directionality: Microwaves pass from port 1 to port 2, from port 2 to port 3, and from port 3 to port 1. The result is the same as with the \(\pm \pi/2\) shifts: the circulators break TRS, resulting in identical GUE spectra for the two disconnected subgraphs, but with an opposite sense of propagation within the respective subgraphs. Again the two subgraphs may thus be described in terms of a secular matrix \(h_0\) and its complex conjugate \(h_0^\ast\), respectively.

The phase difference between the two connecting bonds is adjustable by means of mechanical phase shifters which in reality, however, do not change the phase but the length. This approach has the shortcoming that for a given length change \(\Delta l\) the phase shift \(\Delta \varphi\) depends on frequency \(\nu\):

\[
\Delta \varphi = k \Delta l = \frac{2\pi \nu}{c} \Delta l
\]

where \(k\) is the wave number, and \(c\) is the vacuum velocity of light. \(l = nl_0\) is the optical length where \(l_0\) is the geometrical length and \(n = 1.43\) the index of refraction of the dielectric within the coaxial cables.

Figure 1(b) shows the schematic drawing of one realized graph and (c) the photograph of the corresponding experimental realization. The bonds of the graphs were formed by Huber & Suhner EZ-141 coaxial semi-rigid cables with SMA connectors, coupled by T junctions at the nodes. The phase shifters (ATM, P1507) had been equipped with motors to allow for an automatic stepping. Reflection and transmission measurements were performed with an Agilent 8720ES vector network analyzer (VNA) with the two ports \(P_1\) and \(P_2\) at equivalent positions of the two subgraphs. The corresponding reflection and transmission amplitudes will be denoted in the following by \(S_{ij}\), where \(i,j = 1,2\) is defined by the port. The operating range of the circulators (Aerotex I70-1FFF) positioned at nodes 7 and 7 extended from 6 GHz to 12 GHz. Therefore the analysis of the data was restricted to this window.

We started by taking a series of measurements for constant \(\Delta l\). Figure 2(top) shows the transmission for different \(\Delta \varphi\) are stacked onto each other. (Bottom) The same data, but rearranged to constant \(\Delta \varphi\).
There is a complete equivalence to the Zeeman splitting of spin doublets: In the present experiment the antiunitary symmetry is destroyed when departing from the \( \pi \) line, whereas for conventional spin systems this effect occurs when applying a magnetic field. This is a clear confirmation that we had been successful in constructing a graph with anti-unitary symmetry \( T \) with \( T^2 = -1 \). The distances of the six Kramers doublets seen in Fig. 3 at the \( \pi \) line are equal within 20 percent. This shows a clear tendency of the levels towards an equal level spacings at the \( \pi \) line, one of the fingerprints for a GSE spectrum.

To obtain the complete eigenfrequency spectrum, we proceeded as follows: Though there are two coupled channels, they are equivalent to one effective single channel for \( \Delta \varphi = \pi \) due to symmetry. In this case the scattering matrix reduces to a phase factor \( S = e^{i\varphi} = (1 - iK)/(1 + iK) \), i.e. \( \varphi = -2\arctan(K) \). \( K \) may be written as a sum over resonance poles \( a_n/(x - x_n) \) meaning step-wise phase changes at \( x = x_n \). By taking the phase derivative these steps turn into sharp peaks with widths limited by absorption (which had been discarded in the argumentation). This allowed for an automatic determination of about 90% of the eigenvalues. With the additional information from the spectral level dynamics, see Fig. 3, the missing ones could be easily identified. About 10% of the Kramers doublets split due to experimental imperfections. Whenever this was evident from the level dynamics, the resulting two resonances were replaced by a single one in the middle.

The integrated density of eigenfrequencies may be written as \( n(\nu) = n_{\text{Weyl}}(\nu) + n_{\text{fluc}}(\nu) \) where the average part is given by Weyl’s law, \( n_{\text{Weyl}}(\nu) = (\pi/L)(2\nu/c) \) with \( L \) denoting the sum of all bond lengths \( L \). The fluctuating part \( n_{\text{fluc}}(\nu) \) reflects the influence of the periodic orbits \( \nu \). We determined \( n_{\text{fluc}}(\nu) \) by fitting a straight line to the experimental integrated density of eigenfrequencies and subtracting the linear part. A small number of missing or misidentified resonances showing up in step-wise changes of \( n_{\text{fluc}}(\nu) \) enabled the final correction of the spectrum. From the fit the length was obtained, e.g., \( L = 2.93 \) m for the graph shown in Fig. 1(c). The nearest neighbor spacings \( s \) are calculated as the difference \( s_n = n_{\text{Weyl}}(\nu_{n+1}) - n_{\text{Weyl}}(\nu_n) \) for the individual graph guaranteeing a mean level spacing \( \langle s \rangle = 1 \).

Figure 4 shows the resulting distribution of spacings between neighboring levels in units of the mean level spacing. To improve the statistics, the results from 8 different graphs were superimposed, leading to 1006 Kramers doublets. The red solid and the green dotted line correspond to the Wigner prediction for the GSE,

\[
p_{\text{GSE}}(s) = \frac{2^{18}}{3\pi^3} s^4 \exp\left(-\frac{64}{9\pi^2} s^2 \right)
\]

and the GUE,

\[
p_{\text{GUE}}(s) = \frac{32}{\pi^3} s^2 \exp\left(-\frac{4}{\pi^2} s^2 \right)
\]
respectively. The experimental result fits well to the GSE distribution and, though the statistical evidence as yet is only moderate, it is clearly at odds with a GUE distribution. The inset of Fig. 4 shows the associated spectral rigidity $\Delta_3(L)$ [1]. Again a good agreement with the GSE random matrix prediction is found. The small deviations suggest some percents of misidentified levels, which would have only a minor influence on the nearest neighbor spacings distribution but would distort long-range correlations.

It needed half a century after the establishment of random matrix theory by Wigner, Dyson, Mehta and others to arrive at the present experimental realization of the third of the three classical random matrix ensembles. It might be considered surprising that two bonds between the two subgraphs are already sufficient to turn the two GUE spectra of the disconnected subgraphs into a GSE spectrum for the total graph. On the other hand the GUE spectra of the disconnected subgraphs are already sufficient to turn the two central rigidity $\Delta_3$ into two Gaussian orthogonal ensemble (GOE) and GUE in three classical ensembles are present, namely the GSE for $\Delta_\phi = \pi$, the GOE for $\Delta_\phi = 0$, and the GUE in between. However, the most promising future aspect is undoubtedly that the whole spin $1/2$ physics [27] is now accessible to microwave analogue studies. This work was funded by the Deutsche Forschungsgemeinschaft via the individual grants STO 157/16-1 and KU 1525/3-1. CHJ acknowledges the Leverhulme Trust (ECF-2014-448) for financial support.