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3D stochastic bicontinuous microstructures: Generation, topology and elasticity

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Motivated by recent experimental investigations of the mechanical behavior of nanoporous metal we explore an efficient and robust method for generating 3D representative volume elements (RVEs) with strikingly similar behavior. Our approach adopts Cahn’s method of generating a Gaussian random field by taking a superposition of standing sinusoidal waves of fixed wavelength but random in direction and phase. In its theory part, our study describes closed-form expressions for how the solid volume fraction affects the binarization level, mean structure size, specific surface area, averages of mean and Gaussian curvature, and the scaled topological genus. Based on numerical studies we report on criteria for achieving representative realizations of the structure by proper choice of the number of waves and element size. We also show that periodic structures are readily created. We analyze the mechanical properties considering linear and infinitesimal elasticity and evaluate the residual anisotropy (which can be made small) and the effective values of the Young’s modulus and Poisson’s ratio. The numerical results are in excellent agreement with experimental findings for the variation of stiffness with solid fraction of nanoporous gold made by dealloying. We propose scaling relations that achieve naturally a perfect agreement with the numerical and experimental data. The scaling relation for the stiffness accounts for a percolation-to-cluster transition in the random field microstructure at a finite solid fraction. We propose that this transition is the origin of the previously reported anomalous compliance of nanoporous gold.

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1. Introduction

Stochastic bicontinuous microstructures, in which two contiguous phases interpenetrate, are characteristic for many materials formed by decomposition of a homogeneous parent phase. Examples are spinodally decomposed metal alloys [1,2], polymer blends [3,4] or microemulsions [5], foams that can be polymeric, metallic or food [6] and network solids such as nanoporous metals made by dealloying solid solutions [7,8]. The properties of such microstructures depend on geometric or topological characteristics including the specific surface area, the tortuosity of transport paths through one or both of the phases, the characteristic structure size, and measures for connectivity, such as e.g. genus per unit volume. Here, we describe how spinodal-like stochastic microstructures can be generated by a convenient and fast numerical algorithm and we explore their geometric and mechanical properties.

Our study is motivated by recent research in the field of nanoporous metals made by dealloying. A substantial body of experiments explores specifically the mechanical behavior of nanoporous gold made in this way [9–18]. Nanoporous materials made by dealloying can be understood as networks of nanoscale struts or “ligaments”, typically with solid volume fractions between 0.25 and 0.50 [7–9]. While their microstructure is stochastic, their mechanical behavior is typically discussed with reference to the Gibson-Ashby scaling relations [6] which have been supported - on the modeling side - by periodic structures. However, the mechanical properties of dealloying-made nanoporous solids are not well described by the Gibson-Ashby scaling relations, and specifically their Young’s modulus can be more than an order of magnitude less than predicted [14,19]. While surface excess elasticity has been ruled out as an explanation [20,21], nonlinear elastic behavior of the bulk may contribute [22]. Yet, the decisive issue appears to be
network connectivity [13,14,23,24], which may systematically vary with ligament size and solid fraction.

Cahn [1] showed that the composition field in the early stages of spinodal decomposition could be approximated by the superposition of waves with random direction but same wavelength. In one of the first computer simulations of microstructure evolution in materials science, Cahn showed that binarizing this composition field indeed yielded the bicontinuous microstructure of spinodally decomposed mixtures [1]. In the present work, we explore a modern numerical implementation of Cahn’s algorithm, which efficiently generates 3D stochastic bicontinuous microstructures. We present closed-form expressions for specific surface area and scaled connectivity as a function of the wavelength and solid fraction. Using finite element simulations, we document the elastic behavior and the requirements on the size of the volume element (VE) for achieving representative behavior. The numerical part of our study specifically focuses on the comparison to nanoporous gold, from which the experimental validation of our model emerges. We point out that a percolation-to-cluster transition at finite density can be related to the anomalously low stiffness of experimental nanoporous gold. We bring to the reader’s attention an alternative scaling law, suggested by Roberts and Garboczi [25], which accounts for the transition and remains accurate over a wide range of solid volume fraction. Since many numerical simulation approaches use periodic boundary conditions, we show how volume elements with translational periodicity can be generated. We also provide a comparison between the results for aperiodic and periodic microstructures considering the predicted effective elastic properties and demonstrate that they are in very good agreement.

2. Methods of 3D generation of nanoporous metal microstructures

The 3D geometry of the microstructure of nanoporous gold has been characterized in experiment by tomographic reconstruction. Approaches were based on transmission electron microscopy [26–28], X-ray nanotomography [30–32], focused-ion-beam (FIB) sectioning [23,24,33,34] and atom-probe tomography [35]. These approaches reveal a random network of ligaments which are interconnected in nodes. The surface is found dominated by convex and saddle-shaped patches [33], so that the average of the mean curvature is positive [26,33]. The connectivity density is less than in regular geometric networks such as the gyroid structure [23], and it varies little when the mean structure size is varied by annealing-induced coarsening [24,33]. Experimental 3D microstructures have been used as the basis for numerical simulation of the mechanical response of nanoporous gold [23,24,32]. Studies working with large volume elements, as better representatives for macroscopic material response, find strength and stiffness in qualitative agreement with experimental mechanical tests [23,24].

While experimental reconstructions are distinguished by their realistic geometry, they are also restricted to specific realizations within the family of possible structures with the same stochastic building principles. Furthermore, systematic experiments exploring variations of the microstructure of nanoporous gold with solid fraction remain yet to be reported. Numerical simulation approaches, and specifically atomistic approaches such as molecular dynamics, also often require models that are compatible with 3D translational periodicity. In each instance, the available experimental database for 3D structures is not yet sufficient for comprehensive numerical studies. This highlights the need for the computational generation of model structures.

Computational generation methods often use deterministic periodic unit cell-based idealizations. This may involve heuristic constructions or constructions of (analytical) level surfaces, e.g., triply periodic minimal surface-based unit cells. The Kelvin model [36] consisting of a regular packing of tetra(kaidecahedra, its variations by Waire and Phelan [37] and the model of Gibson and Ashby [6,38] are among the most widely used idealized unit cells for nanoporous metals. Agglomeration of the mass in junctions is considered making use of modified rectangular unit cells in Refs. [39,40]. Periodic diamond cubic unit cells were recently used for modeling nanoporous gold samples [21].

Constructions of triply periodic bicontinuous cubic microdomain morphologies, which are generated by making use of triply periodic continuous minimal surfaces [41,42], constitute another attractive direction. These level surfaces and the developed microstructures have the symmetries of a crystallographic group such as cubic, tetragonal, rhombohedral, and orthorhombic symmetries. Their smooth surfaces allow incorporation of theories accounting for surface curvature effects more realistically. For the use of triply periodic minimal surface based microstructures in modeling of nanoporous gold samples, see e.g. Refs. [23,43,44]. Gyroids as approximants to nanoporous metal foams were investigated in Ref. [45]. Moreover, microstructures made up of single and double gyroids were used in computation of the specific surface area of nanoporous materials in Ref. [46].

Although periodic unit cell idealizations prove efficient and simple, they fall short in reflecting certain key morphological, topological and mechanical characteristics of nanoporous metals. On the morphological part, the 3D reconstructions emphasize a random, as opposed to periodic, structure. So far, no evidence of anisotropy in the mechanical response has been reported for nanoporous gold and the reconstructions appear isotropic. On the contrary, the structures generated by making use of the aforementioned heuristic constructions or constructions of (analytical) level surfaces show cubic anisotropy with high polarity in terms of their directional dependence of Young’s modulus.

It is possible to bridge this gap by considering stochastic disorder in cell structures, see, e.g., [47–49]. Voronoi or Laguerre tessellations constitute a more systematic method in forming stochastic cells [50–53]. A specific step in this direction was the introduction of disorder in the diamond-based network structure of Ref. [21], which indeed provided qualitatively improved agreement with experiments.

Computational generation of stochastic bicontinuous geometries of nanoporous metals are often created at considerable computational expense. With attention to modeling dealloying-made metal network structures, several studies have used the simulation of spinodal decomposition via phase-field [54–56] or kinetic lattice Monte Carlo approaches [19,57]. These approaches achieve a striking similarity to the experimental reconstructions of dealloying-made metal network structures.

A computationally more efficient method of modeling phase-separation dynamics, i.e., the phase-ordering dynamics of thermodynamically unstable phases, is the use of cell dynamical systems [58–60], e.g., coupled maps and cellular automata.

Making use of leveled Gaussian random fields, in which the interfaces between cells are defined by level cuts of random fields [61], in computational generation of stochastic bicontinuous geometries constitutes an even more efficient approach. Leveled-wave models for random fields were investigated for bicontinuous material morphologies formed due to phase-separation [1], have found attention in various applications [25,62–64].
3. Gaussian random fields and microstructure generation

We start by introducing the notion of a continuum with microstructure and computational micro-to-macro transition [65,66] as sketched in Fig. 1. Let \( M \subseteq \mathbb{R}^3 \) denote the homogenized macrocontinuum. A typical point \( M \subseteq \mathbb{R}^3 \) has a representative volume \( V \subseteq \mathbb{R}^3 \) that consists of two phases: the solid bulk phase, i.e., the metallic microstructure, and the pore phase, with their corresponding volumes denoted by \( B \subseteq \mathbb{R}^3 \) and \( P \subseteq \mathbb{R}^3 \), respectively. We then have \( V = B + P \), and the interface of the microstructure is denoted by \( \partial V \).

3.1. A specific random field

The microstructure considered here derives from the random field which is generated by superimposing standing sinusoidal waves of fixed wavelength and amplitude but random direction and phase [1]:

\[
\begin{align*}
\mathbf{f}(\mathbf{x}) &= \frac{1}{\sqrt{N}} \sum_{i=1}^{N} \sin(q_i \cdot \mathbf{x} + \phi_i) \\
C_2(r) &= \langle \mathbf{f}(\mathbf{x}) \cdot \mathbf{f}(\mathbf{x} + \mathbf{r}) \rangle = \frac{\sin(q_0 |r|)}{q_0^2}
\end{align*}
\]

where brackets denote ensemble averages over realizations, and \( r = |x_1 - x_2| \) (see Supplementary Material for the derivation details). If \( N \) is sufficiently large, the value of the random function \( f(\mathbf{x}) \) at a given position \( \mathbf{x} \) follows a Gaussian distribution with \( \mathcal{P}(f) = e^{-f^2/2}/\sqrt{2\pi} \) and the mean wave direction \( \mathbf{v} = \frac{\nabla \mathbf{f}(\mathbf{x})}{|\nabla \mathbf{f}(\mathbf{x})|} \), approaches zero as a consequence of the central limit theorem. Then, \( f \) is an isotropic function, see Supplementary Material.

Given the random function \( f \), the different phases of the system are defined via a level cut \( \xi \):

\[
\begin{align*}
x \in B & \quad \text{if } f(\mathbf{x}) < \xi, \\
x \in \partial B & \quad \text{if } f(\mathbf{x}) = \xi, \\
x \in P & \quad \text{if } f(\mathbf{x}) > \xi.
\end{align*}
\]

In this study, we create numerical realizations of the random field microstructure on regular 3D voxel lattices by summing up \( N \) waves (typically \( N = 10000 \)) with randomly oriented \( \mathbf{q}_i \) and random phases \( \phi_i \). We explore the variation of properties as the structure was repeatedly created with different initial settings of the random number generator; the individual structures in such studies are referred to as “realizations”. Examples of microstructures for different values of \( \xi \) and for volume element size of 12 wavelengths are shown in Fig. 2.

3.2. Generating periodic structures

The microstructures generated using Eq. (1) along with Eq. (3) are generally not periodic. Periodicity can be accomplished by selecting a finite number of waves which have integer wave number in all directions and constant modulus. Our starting point is the field given by Eq. (1). We take \( \mathbf{e}_1 \), \( \mathbf{e}_2 \), \( \mathbf{e}_3 \) unit vectors of an orthonormal basis in real space and let all \( \mathbf{q}_i \) be of the form

\[
\mathbf{q} = \frac{2\pi}{a} (h, k, l)
\]

where the Miller indices \( h, k, l \) are integers and \( a \) is a constant. Then, \( f(\mathbf{x}) = f(\mathbf{x} + m\mathbf{e}_1 + n\mathbf{e}_2 + o\mathbf{e}_3) \) if \( m, n, o \) are arbitrary integers. This implies that \( f \) has translational periodicity with lattice vectors of magnitude \( a \). Furthermore, if all \( \phi_i \) in Eq. (1) are identical, then \( f \) is also invariant with respect to exchange of the axis and so \( f \) has the symmetry of a cubic lattice with lattice parameter \( a \).

The above considerations imply that “spinodal” structures with translational periodicity or even cubic symmetry may be constructed by restricting the sum in Eq. (1) to a set of \( \mathbf{q}_i \) with a given, constant value \( H = \sqrt{h^2 + k^2 + l^2} \), so that \( |\mathbf{q}| = q_0 = 2\pi h/a \) is a constant. The function \( f \) then forms a periodic repetition of identical cubic unit cells with edge length \( a \).

Real spinodal structures emerge by the growth of a continuum with initial fluctuations including random orientation of their wave vectors [1]. It is then expected that unit cells containing many waves are more representative of real structures than those that contain few waves, which in turn suggests the use of values of \( H \) that maximize the number of individual vectors \( (h, k, l) \). In diffraction theory, this corresponds to Bragg reflections with high multiplicity. For instance, \( H = \sqrt{160} \) (which implies \( M = 10.3 \) elements per edge length, see Section 4.2 below) is only compatible with \((1240)\)-type vectors, which only gives 12 independent directions. By contrast, the marginally larger \( H = \sqrt{161} \) is compatible with \((1241)\), \((1162)\), \((1065)\) and \((984)\)-type vectors and, thus, 96 independent directions. Similarly, \( H = \sqrt{146} \) is compatible with \((0511)\), \((1112)\), \((189)\), \((3411)\), and \((479)\)-type vectors, leading also to 96 independent directions.

Fig. 3 illustrates the wave vector directions for \( H = \sqrt{146} \), along with binarized structures created from these waves in Eq. (1). The structure in Fig. 3 (c) is based on setting \( \phi_i = 0 \) and exhibits full cubic symmetry. It is striking that ring-like features near the center of each face have no equivalent in experimental stochastic structures. This problem is solved by using random phase shifts \( \phi_i \) in Eq. (1), see Fig. 3(b). The resulting structure compares well to experimental microstructures such as those of Dealloying-made nanoporous gold. The structure now lacks the rotation- and mirror symmetry of a cubic lattice, but it retains the full translational periodicity and is suited for periodic boundary conditions.
4. Morphological characterization

In this section, we present explicit analytical results that characterize the morphological and topological properties of the stochastic aperiodic microstructures generated with Eq. (1) in the limit of a large number $N$ of waves.

4.1. Solid volume fraction

For porous materials, the volume fraction of phases constitutes one of the most important morphological descriptors. Using $|\mathcal{V}| = |\mathcal{B}| + |\mathcal{P}|$, with $|\{ \bullet \}|$ denoting the volume contained in $\{ \bullet \}$, the volume fractions, $\phi_{\mathcal{B}}$ and $\phi_{\mathcal{P}}$, of the solid and pores, respectively, are defined as

$$
\phi_{\mathcal{B}} := \frac{|\mathcal{B}|}{|\mathcal{V}|} \quad \text{and} \quad \phi_{\mathcal{P}} := \frac{|\mathcal{P}|}{|\mathcal{V}|}.
$$

(5)

with $\phi_{\mathcal{B}} + \phi_{\mathcal{P}} = 1$. Given the Gaussian properties of the random field (1) the volume fraction $\phi_{\mathcal{B}}$ of the microstructure defined by Eq. (3) is given by

$$
\phi_{\mathcal{B}} = \frac{1}{2} \left[ 1 + \text{erf} \left( \frac{\xi}{\sqrt{2}} \right) \right].
$$

(6)

where $\text{erf}(x)$ denotes the error function. Inverting the equation above, we can have any desired volume fraction by setting the level cut $\xi$ to

$$
\xi(\phi_{\mathcal{B}}) = \sqrt{2}\text{erf}^{-1}(2\phi_{\mathcal{B}} - 1),
$$

(7)

where $\text{erf}^{-1}(x)$ denotes the inverse error function. Fig. 2 shows microstructures with solid volume fractions from 0.10 to 0.50, where we can see how the microstructure with 0.10 solid volume fraction is formed of many disjoint regions. This is consistent with the percolation threshold $\phi_{\mathcal{B}}^p = 0.159$ (see below). This level of disjointness does not allow the structure to withstand applied loads.

4.2. Characteristic structure size and mean ligament diameter

Each of the components $\cos(\mathbf{q}, \mathbf{x})$ that make up the random field $f(\mathbf{x})$ in Eq. (1) has the identical wavelength $\lambda = 2\pi/q_0$, which could be used as a first guess at a characteristic structure size. However, we note that a view that accounts more specifically for the properties of the random field $f(\mathbf{x})$ is to inspect the autocorrelation function, Eq. (2). The first maximum of that equation is at $q_0f = 1.23 \times 2\pi$, which measures a characteristic distance, $l$, between regions of maximum $f$. In the binarized structures, we can interpret $l$ to represent the characteristic structure size as measured by the mean distance between local centers of the solid or the pore space. We, thus, have
with \( a = 1.23 \). Experiments with dealloying-made nanoporous structures often report a characteristic size \( L \) representing a mean diameter of the solid struts or ligaments that can be perceived in medium density binarized structures such as those of Fig. 2. For structures with equal volume fraction of the two phases, symmetry suggests that \( L = L/2 \).

For periodic microstructures, in view of Eq. (8), we can take the characteristic structure size as \( L = 1.23 \times 2\pi/\rho_0 = 1.23a/\rho \). This means there are \( M = aL = H/1.23 \) microstructural elements per edge length of the unit cell, and the unit cell volume \( a^3 \) contains about 0.54H \( L \) microstructural elements.

We note that another important quantity that can be used to characterize the properties of the two-phase system is the auto-correlation function, \( \Gamma_2(r) \), of the solid phase. \( \Gamma_2(r) \) is defined as

\[
\Gamma_2(r) = \langle I(x_1)I(x_2) \rangle,
\]

where \( I(x) \) is the indicator function for the solid phase which is defined via the random function as

\[
I(x) = \begin{cases} 1, & \text{if } x \in B, \\ 0, & \text{otherwise}. \end{cases}
\]

Furthermore, \( r = |x_1 - x_2| \). The form of the autocorrelation function \( \Gamma_2(r) \) reflects certain features of the microstructure, e.g. short range order and spatial correlation [61]. Fig. 4(a) shows plots of the numerically computed \( \Gamma_2 \) normalized autocorrelation function \( \Gamma_2^{\text{norm}}(r) = \langle \Gamma_2(r) - \phi_B^2 \rangle / (\phi_B - \phi_B^2) \) for volume fractions from 0.20 to 0.50. As we can see, all plots exhibit oscillations of identical periodicity for small \( r \), hence signaling a short range order and spatial correlation in the medium generated. As expected, the mean distance between successive peaks of \( \Gamma_2^{\text{norm}}(r) \) is approximately \( \lambda = 2\pi/\rho_0 \), which is the same for \( C_2(r) \), given by Eq. (2).

### 4.3. Surface area-to-volume ratio

The ratio of surface area to total (solid plus pores) volume is

\[
S := \frac{|\partial B|}{|V|},
\]

where \( |\partial B| \) denotes the total area associated with \( \partial B \). By considering the behavior of the two-point correlation function, \( C_2(r) \), near \( r = 0 \), and making use of results from the scattering theory, we find an explicit formula for the surface-to-volume ratio as (see Supplementary Material for details)

\[
S = \frac{2\rho_0}{\pi^3} e^{-\ell^2/2}.
\]

The ratio of surface area to solid volume is then

\[
S_B = \frac{1}{\phi_B} \frac{2\rho_0}{\pi^3} e^{-\ell^2/2}.
\]

The dimensions of both \( S_B \) and \( S \) are \( [S] = [S_B] = [L] \). Thus, multiplication with the characteristic size \( L \) supplies non-dimensionalization and provides quantities devoid of volume and length scale differences. Fig. 4(b) shows the plots of \( S_B(\phi_B) \) and \( S(\phi_B) \) as given by the above relations, where we can see, as expected, that \( S_B \) decreases monotonically, and \( S \) is symmetric with respect to \( \phi_B = 0.50 \) with a maximum at this point. The discrete points in Fig. 4(b) correspond to the numerical values computed from the auto-correlation function \( \Gamma_2(r) \) of the generated microstructures (see the inset of Fig. 4(a) and the Supplementary Material for details on how \( \Gamma_2(r) \) is related to \( S \), observing an excellent agreement with the analytical predictions).

### 4.4. Average curvatures and percolation-to-cluster transition

In differential geometry, the product of principal curvatures \( \kappa_1 \) and \( \kappa_2 \) is given by the Gaussian curvature \( K \) which for planar surfaces is related to the area-to-volume ratio \( S \) via the equation

\[
K = \frac{S}{h^2},
\]

where \( h \) is the mean thickness of the structure. For microstructurally random materials, the mean thickness is related to the mean free path \( \ell \) via the equation

\[
h = \frac{\ell}{\phi_B}.
\]

Thus, the Gaussian curvature of the solid phase is given by

\[
K = \frac{S_B}{\phi_B^2},
\]

where \( S_B \) is the surface-to-volume ratio of the solid phase. The average Gaussian curvature of a random microstructure is given by

\[
\langle K \rangle = \int_0^1 \langle K \rangle(\phi_B) \phi_B d\phi_B,
\]

where \( \langle K \rangle(\phi_B) \) is the average Gaussian curvature at volume fraction \( \phi_B \). This quantity can be used to study the percolation-to-cluster transition, which occurs when \( \langle K \rangle \) reaches a critical value, \( K_c \), at which point the microstructure undergoes a phase transition from a disordered to a clustered state. For example, if \( K_c = 1 \), then the percolation threshold is \( \phi_B^c = 0.50 \). The critical values of \( \langle K \rangle \) can be determined numerically using Monte Carlo simulations or other numerical methods.
and $\kappa_2$ yields the Gaussian curvature of a surface at a point as $\kappa_g = \kappa_1\kappa_2$, whereas the mean curvature is defined as $\kappa_m = [\kappa_1 + \kappa_2]/2$. For the Gaussian random field of Eq. (1), the average Gaussian curvature and the average mean curvature are [64]

$$\langle \kappa_g \rangle = \frac{q_0^2}{6} \left[ \frac{x^2}{2} - 1 \right] \quad \text{and} \quad \langle \kappa_m \rangle = -\xi q_0 \frac{1}{2} \sqrt{\frac{\pi}{6}}$$

(14)

respectively, where we use the curvature definition of $\kappa_1 > 0$ and $\kappa_2 > 0$ for a solid (convex) ellipsoid. Table 1 shows tabulated values of surface area-to-volume ratios, average Gaussian and average mean curvature, rescaled with the characteristic structure size $L$ given by Eq. (8) as a function of the solid volume fraction. Fig. 5(a) shows the variation of average scaled Gaussian and mean curvatures with the solid volume fraction, along with examples of microstructures.

We observe that the Gaussian curvature crosses zero at two critical points, namely $\phi_0 = \phi_{g0}^P = 0.159$ and $\phi_0 = 1 - \phi_{g0}^P = 0.841$, which correspond to the solid percolation threshold and pore percolation threshold, respectively. Based on these values, we can identify four distinct types of microstructure: (i) $\phi_0 < \phi_{g0}^P$, the microstructure is formed of disjoint solid phases (convex ellipsoids) with local curvatures $\kappa_1 > 0$ and $\kappa_2 > 0$ which gives $\kappa_m > 0$ and $\kappa_0 > 0$, (ii) $\phi_0 < \phi_{g0}^P < 0.50$, thin and long saddle-like ligament formations with sharp positive local curvature ($\kappa_1 > 0$) at their sections but mild negative local curvature ($\kappa_2 < 0$) along their lengths dominate with $|\kappa_1| > |\kappa_2|$ resulting in $\kappa_m > 0$ and $\kappa_0 < 0$, (iii) $0.50 < \phi_0 < 1 - \phi_{g0}^P$, thick and short saddle-like ligament formations with mild positive local curvature ($\kappa_1 > 0$) at their sections but sharp negative local curvature ($\kappa_2 < 0$) along their lengths dominate with $|\kappa_1| < |\kappa_2|$ resulting in $\kappa_m < 0$ and $\kappa_0 < 0$, (iv) $\phi_0 > 1 - \phi_{g0}^P$, disjoint pores (concave ellipsoids) characterize the microstructure with $\kappa_1 < 0$ and $\kappa_2 < 0$, and so $\kappa_m < 0$ and $\kappa_0 > 0$.

We also note that only the 0.50 vol fraction case has zero average mean curvature, which agrees well with the results reported in Ref. [68]. At this point the average Gaussian curvature has a minimum. The Gaussian curvature is symmetric with respect to $\phi_0 = 0.50$, whereas the mean curvature is instead showing a monotonic decrease with increasing phase volume fraction. These trends are in agreement with experimental studies presented in the literature. Fujita et al. [27] measured near zero mean curvature for nanoporous gold samples of $\phi_0 = 0.50$. A positive average mean curvature was reported by Rosner et al. [26] for nanoporous gold with a lower metal volume fraction of $\phi_{g0} = 0.24$.

### 4.5. Genus

The genus $G$ of a surface describes its connectivity; it can be understood as the number of continuous tunnels or loops in the structure. For instance, a sphere, donut and pretzel have the genus 0, 1, and 2, respectively. By making use of the Gauss-Bonnet theorem and the known Gaussian curvature per unit volume, it is possible to obtain the following expression for the genus per unit volume, $G_v$, in our random field microstructures (see Supplementary Material for details):

$$G_v = \frac{1}{12\pi^2 \sqrt{3}} \left[ 1 - \frac{x^2}{2} \right] e^{-c^2/2}$$

(15)

The above relation depends on the characteristic length scale of the microstructure through $q_0$. To study a non-dimensional quantity we rescale the genus per volume with the characteristic length scale $L$, obtaining the normalized genus per unit volume $G_v$. The size $L$ is the equilultral size of the spherical domain occupied by the phase of interest.
structure size L, see Eq. (8), to obtain the scaled genus per volume:
\[ g_s = G_v \times \frac{L^3}{3} \left( 1 - \varepsilon^2 \right) e^{-\varepsilon^2/2}. \] (16)

Fig. 5(b) shows the prediction of Eq. (16) for the scaled genus versus solid fraction. Similarly to the average curvatures analysis from the previous section, Eqs. (15) and (16) also predict a topological transition at a finite phase fraction. The genus vanishes at \( g = 1 \), which according to Eq. (6), corresponds to \( \phi_p = \phi_p^0 = 0.159 \) for \( g = 1 \). For phase fractions \( \phi_p > \phi_p^0 \), the structure may be understood as a percolating and multiply interconnected network. At lesser \( \phi_p^0 \) connectivity is lost and the structure becomes an array of isolated clusters with genus zero. The phase fraction \( \phi_p^0 \) thus marks the percolation-to-cluster transition, below which Eq. (16) is no longer applicable. On the other hand, at \( \phi_p = 0.841 \), which corresponds to \( \varepsilon = 1 \), there is the topological transition above which the structure contains isolated voids.

To compute the genus numerically we use the fact that the genus is related to the second Betti number as \( B_2 = \frac{1}{2} g \). Thus, an assessment whether numerical realizations are representative can be based on their degree of mechanical anisotropy.6

Supplementary Material.

5.2. Quantification of the degree of anisotropy and aggregate properties

In its limiting behavior for a large number \( N \) of waves and a large volume, the Gaussian random field of our analysis is isotropic, and for the current particular case this property implies elastic isotropy.6 Thus, an assessment whether numerical realizations are representative can be based on their degree of mechanical anisotropy. An interpretation of the elasticity constants is in terms of Young’s modulus \( E \), i.e., the sensitivity of stress increment to strain increment along the loading direction under uniaxial state of stress. Assuming loading along the unit direction \( m \) the corresponding effective Young’s modulus \( E_m^* \) is computed as
\[ E_m^* = \frac{1}{[m \otimes m] : \mathbf{S} : [m \otimes m]} \] (17)

where \( \mathbf{S} = [\mathbf{C}^*]^{-1} \) is the compliance tensor.7 By plotting \( E_m^* \) for all \( m \) one can develop an understanding regarding the directional dependence of elasticity.

For the quantification of the extent of anisotropy in the material, various measures exist, e.g., for cubic crystals, Zener [71] introduced an anisotropy index. However this index is restricted to cubic crystals and lacks universality [72]. Thus, we use the universal anisotropy index [72], which overcomes these limitations:
\[ A^V = \frac{K_V}{K_R} + \frac{5}{3} \frac{\mu_V}{\mu_R} - 6 \geq 0. \] (19)

Here, \( K \) and \( \mu \) represent isotropized single crystal elasticity shear and bulk moduli which are obtained through assignment of the single crystal orientation uniformly on a sphere of 3D distributions. The subsequent ensemble averaging over the unit sphere6 is denoted by \( \langle \cdot \rangle \), so that the elastic constitutive tensor and the compliance tensor are given as
\[ \mathbf{C}_V^* = \langle \mathbf{C}^* \rangle = 3K_V^{vol} + 2\mu_V^{dev}, \]
\[ \mathbf{S}_R^* = \langle \mathbf{S}^* \rangle = \frac{1}{3K_R^{vol}} + \frac{1}{2\mu_R^{dev}}. \] (22)

where \( R \) and \( V \) stand for Reuss and Voigt estimates, respectively. \( A^V = 0 \) for isotropic materials, as in the case of macroscopic

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4 For the details of how these reported results are converted to our scaling convention, the reader is referred to Appendix A.
5 Due to the porous structure \( \mathbf{M}_e = 1/|\Omega| \int_{\Omega} dV \). Thus, macroscopic strains are computed through control nodes of the volume element.
6 Statistical isotropy does not necessarily result in mechanical isotropy [61].
7 Poisson’s ratio, where the direction of observation \( m_i \) is perpendicular to the loading direction \( m \) can be given as [70].
8 Spherical averaging reads.
\[ \langle \cdot \rangle = \frac{1}{4\pi} \int_0^{2\pi} \int_0^{\pi} \int_0^{2\pi} \langle \cdot \rangle |\Theta| \sin \Phi \sin \Theta d\Phi d\Theta d\Phi. \] (20)

Numerically, this can be realized considering \( M > 200 \) points as unit directors of a Fibonacci tiling, with \( \langle \cdot \rangle_{i} = \langle \cdot \rangle |\Theta_i| \Phi_i \) where \( i = 1, 2, 3, \ldots M \). Then
\[ \langle \cdot \rangle = \frac{1}{M} \sum_{i=1}^{M} \langle \cdot \rangle_{i}. \] (21)

As it is done in this study, for computation of elastic aggregate properties, one can also make use of the invariance properties of the elastic constitutive tensor [73], see Supplementary Material.
response of nanoporous gold, and \( A_{\text{ij}} > 0 \) increases with increasing level of anisotropy\(^9\).

### 6. Results and discussion

#### 6.1. On the topological properties

As pointed out in Ref. [23], microstructures that appear similar to that of nanoporous gold can exhibit variations in apparent scaled genus by one order of magnitude. The observation that the scaled genus of our random field microstructure matches experimental reconstructions of nanoporous gold to the factor two or better may thus be considered as promising. It is also significant that the experimental structures have lesser \( g_c \) than our model. This is consistent with the trend — pointed out by Mameka et al. in Ref. [13] — for \( g_c \) in spinodal-like structures such as nanoporous gold to decrease during spontaneous coarsening. That trend implies the early-stage spinodal structure of Ref [1], which underlies gold to decrease during spontaneous coarsening. That trend implies the early-stage spinodal structure of Ref [1], which underlies gold to decrease during spontaneous coarsening.

According to Ref [13], lesser \( g_c \) may specifically be expected in experiments when samples were allowed to undergo spontaneous coarsening during corrosion, during storage at room temperature, or during deliberate annealing treatments. It is this reasoning that prompts us to limit the comparison between experiment and the reconstructions of nanoporous gold to the factor two or better may eventually involves coarsening and, hence, reduction in \( g_c \).

Motivated by experiments with nanoporous gold made by dealloying (see below), we consider \( \phi_0 \) in the range of 0.20–0.50. The crystal lattice of gold has cubic symmetry and is elastically anisotropic. In the interest of conciseness, this elastic anisotropy is ignored in the current study. We assume isotropic elasticity of the solid phase and, for contact with the experiments, use the parameters of polycrystalline gold for the Young’s modulus \( E_B \) and for the Poisson’s ratio \( v_B \) of the solid phase: \( E_B = 79 \) GPa and \( v_B = 0.44 \) [82].

The macroscopic Young’s moduli and Poisson’s ratios are now discussed. Since the macroscopic elastic isotropy is not a priori satisfied for all volume element sizes, the computations are presented in terms of aggregate macroscopic elastic properties \( E_{\text{VBR}} \) (top) and \( v_{\text{VBR}} \) as depicted in Fig. 7 for 0.20 and 0.50 phase volume fractions. Aggregate elastic properties are derived from Voigt and

\[ E_{\text{VBR}} = \frac{3E}{2(3v + 1)}, \quad v_{\text{VBR}} = \frac{3v}{2(3v + 1)} \]

which is six if the crystal is locally isotropic [72].

 simulations considering various volume element sizes, see Fig. 6, and computed the corresponding macroscopic properties. In order to collect sufficient statistical information, 15 realizations were generated for each volume element size.

Unless otherwise stated, the results are for aperiodic structures. The homogenization scheme is computationally implemented with the finite element method. As detailed in Section 2, we use a fixed\(^11\) wave number \( q_\text{ref} = q_0 \) and we present the volume element sizes in terms of number of wavelengths. We use 3D full integration first order trilinear finite elements through voxelization \(^7\) of the domain. All reported results use fixed sized voxels, where the voxel size is determined through a mesh convergence analysis, considering the computational memory and time requirements of the structures with higher solid volume fractions within the statistical computations. Since no size dependent constitutive phenomena is used during modeling, the computed mechanical properties are equally valid for any ligament size. The underlying statistical properties require the number of waves approaching infinity. As explained earlier \( N = 10000 \) waves are sufficient in our applications considering isotropy of the field. Thus, for all applications with aperiodic structures we use \( N = 10000 \).

Motivated by experiments with nanoporous gold made by dealloying (see below), we consider \( \phi_0 \) in the range of 0.20–0.50. The crystal lattice of gold has cubic symmetry and is elastically anisotropic. In the interest of conciseness, this elastic anisotropy is ignored in the current study. We assume isotropic elasticity of the solid phase and, for contact with the experiments, use the parameters of polycrystalline gold for the Young’s modulus \( E_B \) and for the Poisson’s ratio \( v_B \) of the solid phase: \( E_B = 79 \) GPa and \( v_B = 0.44 \) [82].

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\(^9\) The origin of the index then stems from the contraction

\[ C_j^e = \left( \frac{V}{S^e} \right) \cdot \left( \frac{\langle S_j^e \rangle}{\langle S^e \rangle} \right) - \frac{K_v}{K_v} \cdot \frac{\langle \rho_v \rangle}{\rho_v} \]

which is six if the crystal is locally isotropic [72].

\(^10\) For a fixed wave number, the microstructures are not triply periodic unless the periodicity condition is explicitly imposed, see Section 3.2. In our study, we mainly consider the more general aperiodic ones. Still, we apply periodic boundary conditions due to their superior convergence properties. This is especially important for multiphase composites with highly contrasted phase mechanical properties for which the kinematic uniform boundary conditions overestimates the effective properties at reasonable volume element sizes. For more discussions on the use of periodic boundary conditions for aperiodic microstructures in homogenization the reader is referred to, e.g., [25,75–80].

\(^11\) Voxel based approaches are problematic for physical problems dependent on surface area (such as surface elasticity, thermal transport, electromagnetic radiation and chemical diffusion) as the surface area is appreciably overestimated [81]. This overestimation cannot be remedied by mesh refinement (e.g., for a sphere meshed as cubic voxels as the element side length over radius of sphere ratio tends to zero, the ratio of the voxel mesh surface area to that of a perfect sphere will tend to \( 3/\sqrt{3} = 1.73 \) - in other words, in the limit, the surface area is overestimated by more than 0.70). As we do not study surface area dependent problems, the voxel mesh is fully sufficient for the purpose of this work.
Poisson’s ratio for 0.50 solid volume fraction. Data is presented as mean value with corresponding wavelengths and effective aggregate properties. The effective properties are judged for the volume element size of 12 wavelengths. Therefore, 12 wavelengths is identified as the RVE size. At this point, $E_{\text{HVR}} = 0.0644 \pm 0.0159$ GPa and $\nu_{\text{HVR}} = 0.1741 \pm 0.0145$ for 0.20 volume fraction, and $E_{\text{HVR}} = 14.322 \pm 0.179$ GPa and $\nu_{\text{HVR}} = 0.2813 \pm 0.0009$ for 0.50 solid volume fraction. The scatter in the results for 0.50 solid volume fraction is smaller than at 0.20. Most remarkably, the Young’s modulus for the former is nearly two orders of magnitude larger than for the latter.

For sufficiently large volume element size, the elastic behavior in the computations approaches the isotropy of the random field. We verify this in 1080 simulations whose results are demonstrated in Figs. 8 and 9. Since $A^1 = 0$ represents isotropy and $A^1 \geq 0$, one anticipates decreasing universal anisotropy index with increasing volume element size. As depicted in Fig. 8, for both volume fractions, our results fall below 0.5 and, thus, sufficiently meet isotropy expectation if the volume element size is 12 wavelengths or larger. For 0.20 and 0.50 solid volume fractions, the anisotropy indices read $0.2270 \pm 0.0890$ and $0.0221 \pm 0.0077$, respectively, where the latter structure has about ten times less anisotropy. Fig. 9 demonstrates that the volume element size influence on the directional dependence of normalized Young’s modulus for 0.20 and 0.50 solid volume fractions. Here, the results are ordered vertically from minimum (bottom) to maximum anisotropy index (top). Plots of the first column of Fig. 9(a), that is for the volume element size of 3 wavelengths, signal a substantial deviation from isotropy. Here, the load transfer along directions where there is lack of percolation of the matter is highly disturbed. Thus, the stiffness along such directions is computed to be quite small and this creates an extreme polarity in the directional response of material. As the volume element size gets larger, the surface morphologies tend to that of a sphere. For 0.50 solid volume fraction, this tendency is much faster. Even for the smallest volume element size, no extreme polarization is observed. At this solid volume fraction, the solid and pore phases are symmetric in content, and percolation prevails.

In addition to the stochastic aperiodic structures, we also considered periodic structures as described in Section 3.2. The effective elastic properties of periodic, but not cubic symmetric, structures were computed based on 15 realizations satisfying the identical behavior as our aperiodic ones with 10000 waves. The effective values with growing volume element size with a narrower error bar. The aggregate effective Young’s modulus $E_{\text{HVR}}$ and aggregate effective Poisson’s ratio $\nu_{\text{HVR}}$ are well converged as the volume element size reaches 12 wavelengths. Therefore, 12 wavelengths is identified as the RVE size. At this point, $E_{\text{HVR}} = 0.0644 \pm 0.0159$ GPa and $\nu_{\text{HVR}} = 0.1741 \pm 0.0145$ for 0.20 volume fraction, and $E_{\text{HVR}} = 14.322 \pm 0.179$ GPa and $\nu_{\text{HVR}} = 0.2813 \pm 0.0009$ for 0.50 solid volume fraction. The scatter in the results for 0.50 solid volume fraction is smaller than at 0.20. Most remarkably, the Young’s modulus for the former is nearly two orders of magnitude larger than for the latter.

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6.3. Scaling law for elastic properties

As demonstrated in the preceding section, by using a RVE size of 12 wavelengths or higher, one can sufficiently reduce the gap between Voigt and Reuss aggregate property bounds and the universal anisotropy index. Therefore, aggregate effective elastic properties in fact represent effective elastic properties of the medium. This allows us to drop the subscript HVR and concisely use $E^*$ and $\nu^*$ while denoting the effective Young’s modulus and Poisson’s ratio, respectively. The variation of the converged effective Young’s modulus results from our study with solid fraction is displayed by the circles in Fig. 10(a). The mean values are obtained using computations conducted over 15 realizations. The first notable observation is that our periodic structures with 96 waves give sensibly the identical behavior as our aperiodic ones with 10000 waves. The mutual consistency supports the validity of both approaches.
stochastic microstructures, Roberts and Garboczi [25] propose a scaling law that explicitly accounts for the presence of a percolation threshold given by

\[ E^* = E_B(\phi_B - \phi_0) / (1 - \phi_0)^m. \]

This motivates our plot of the normalized \( E^* \) versus \( \phi_B / (1 - \phi_B) \) in Fig. 10(b). As derived in Section 4.4, \( \phi_B = [1 + \text{erf}(2^{-1/2})] / 2 \approx 0.159 \) is the percolation threshold of our random field microstructures. Since the linear regression does not extrapolate to \( E^* = E_B \) at \( \phi_B = 1 \), we are led to propose a modified version of the Roberts-Garboczi scaling law in the form

\[ E^* = C_2 \left[ \frac{\phi_B - \phi_0}{1 - \phi_B} \right]^m. \]  

(25)

If averaged over the data for periodic and aperiodic structures, the linear regression suggests \( C_2 = 2.03 \pm 0.16 \) and \( m = 2.56 \pm 0.04 \). The excellently linear behavior of our effective Young’s moduli in the representation of Fig. 10(b) provides convincing evidence for the stiffness to vanish along with the topological genus at precisely the percolation limit. This differs from the statement by Roberts and Garboczi [25] who, exploring structures in a more restricted range of phase fractions and working with their original scaling equation, suggested that \( E^* \) may vanish at a solid fraction which is above the percolation limit [25].

The fit to our numerical results with Eq. (25) is reproduced in Fig. 10(a) as the solid line. Also shown (noncircular symbols) is experimental data from Refs. [13–15,83] for the reduced effective Young’s modulus of the material that motivates our study, nanoporous gold, with different solid fraction. That data underlines the finding – highlighted in the experimental studies and confirmed by atomistic simulation [19,55,56] of nanoporous gold – that the material is generally more compliant than predicted by the Gibson-Ashby law. Our compilation of the experimental data emphasizes that this deviation systematically increases as the solid fraction is reduced. Most remarkable, however, is the excellent agreement of our simulations with experiments.

The variation of the converged effective Poisson’s ratio results from our study with solid fraction is displayed by the circles in Fig. 11. Here, the noncircular symbols show experimental data for the elastic Poisson’s ratio of nanoporous gold from Refs. [15,16]. It is seen that our results agree, within error bars, with the experiment. The agreement confirms the notion that our random field microstructure provides a faithful model for the structure of dealloying-mad nanoporous gold. In comparison with the Young’s modulus, \( \nu^* \) exhibits a more moderate variation with \( \phi_B \). We found a highly linear behavior when plotting \( \nu^* \) versus \log(\phi_B) as a solid line in Fig. 11, which fits to the empirical law

\[ \nu^* = D_1 \log \phi_B + D_2, \]  

(26)

taking \( D_1 \) and \( D_2 \) as fitting parameters. The slopes for the periodic and aperiodic structures are within the range \( D_1 = 0.116 \pm 0.003 \), whereas the extrapolated value for \( \phi_B = 1 \) is \( D_2 = 0.363 \pm 0.002 \). Extrapolating this behavior to the percolation threshold indicates that the Poisson’s ratio may remain finite there, contrary to Young’s modulus. The dashed line in the figure represents the Gibson-Ashby model calculation for the effective Poisson’s ratio. This calculation of a constant Poisson’s ratio of 0.30 falls short to reflect the experimentally observed reduction in the transverse strain to axial strain ratio.

Our findings have two important implications:
Fig. 9. Directional dependence of normalized Young’s modulus for increasing volume element size for (a) 0.20 and (b) 0.50 solid volume fraction. 15 realizations for each volume element size are shown. The results are vertically ordered from less (top) to more (bottom) isotropic. In accordance with Fig. 8, for smaller volume element sizes a high directional dependence of Young’s modulus is observed. A highly reduced directional dependence of Young’s moduli is observed for 0.50 solid volume fraction as compared to the 0.20 solid volume fraction. The colors correspond to the radius at each surface point with red denoting a maximum and blue denoting a minimum which are individually computed for each figure. Elastic isotropy is represented by a sphere. (For interpretation of the references to colour in this figure legend, the reader is referred to the web version of this article.)
Firstly, the agreement suggests that the random-field microstructure of our study is in fact a rather accurate representation of the microstructure in experimental nanoporous gold. This provides strong support for the use of similar structures in atomistic simulation studies of nanoporous gold, such as references [19,55,56].

Secondly, our findings suggest that reducing the solid fraction of nanoporous gold systematically brings its microstructure closer to the percolation-to-cluster transition by reducing its topological genus. That notion is anticipated in Ref. [13] based on findings for the microstructural evolution during coarsening of spinodal-like structures [84,85]. The important role of connectivity for the mechanical behavior of nanoporous gold is also advertized in studies determining measures for the scaled genus of nanoporous gold experimentally [23,24]. While such studies remain to be combined with systematic measurement of the mechanical behavior, convincing support comes from a comparative analysis of stiffness and strength, which invokes connectivity to reconcile the observations [14]. By combining information on the topology and on the effective elastic behavior in a wider range of solid fractions, our observations provide the most direct evidence so far for a link between connectivity and stiffness of nanoporous gold. Specifically, we advertize the impact of a systematic drop in connectivity at lesser solid fraction in the material.

The deviation between the stiffness of nanoporous gold and the Gibson-Ashby scaling has led to the proposition of modified scaling laws. Sun et al. [55] proposed a correction that accounts for a transition between bending and tensile deformation as a function of the solid fraction. Huber et al. [47], in their analysis of scaling, also acknowledged that transition but pointed further to the role of structural disorder. Yet, scaling equations that account for the percolation-to-cluster transition have not previously been discussed in context with nanoporous gold. Our results provide strong evidence that this latter phenomenon may hold the key for understanding the systematic deviation between experimental nanoporous gold stiffness data and the classic foam scaling relations.

7. Conclusions

Based on Cahn’s method of levelled Gaussian random fields made of random superposed standing waves with constant wavelength, we proposed an efficient and robust method for generating 3D microstructures for nanoporous gold. The method allows a priori fulfillment of the desired solid volume fraction by its inherent Gaussian property. Morphological, topological and elastic
properties of the generated structures were investigated. Comparison of the findings with experimental observations proved that the method accurately produces many morphological and topological features of dealloying-made nanoporous gold. We proposed scaling relations for the variation of the Young’s modulus and Poisson’s ratio with the solid volume fraction. These relations qualitatively differ from the Gibson-Asby scaling relations for open-cell foams; the key difference is that systematic changes in the topological genus, which include a percolation-to-cluster transition at a finite solid fraction, are accounted for.

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Appendix A. Conversion between non-dimensional quantities

Specific surface areas $S_B$ and $S$ as well as genus density $G_i$ are dimensional properties, thus, they depend on the selected length scale. Multiplication with a characteristic size supplies non-dimensionalization and provides quantities which no longer depend on system size or length-scale. Different length scales have been used in the literature, for example Refs. [23,86] used the inverse of the surface area per unit volume, which is defined in Eq. (13); Ref. [24] used the mean ligament diameter instead.

In our computations we have used the length $L$ as the characteristic size. As we have mentioned in Section 4.2, for the symmetric case we can relate the mean ligament diameter with the characteristic size. As we have mentioned in Section 4.2, for the symmetric case we can relate the mean ligament diameter with the characteristic size $L = L/2$. However, in general we expect that the mean ligament diameter is given as $L = h(\phi_B)L$, where $h(\phi_B)$ is a function of the solid volume fraction. Numerically we can find this function by computing the mean ligament diameter of the microstructures for different values of $\phi_B$ (see Supplementary Material for details). Fig. 12 shows the numerical value of $L$ for different values of the volume fraction, where we can fit a simple linear behavior with parameters: $L/L = [0.53\phi_B + 0.41]$.

Noting that $\lambda = L/a$, we hence arrive at the following empirical relation:

$$L(L, \phi_B) = [L/\alpha][0.53\phi_B + 0.41].$$

(27)

which we can also use to express the inverse of the surface-to-volume ratio as:

$$S^{-1}(L, \phi_B) = \frac{\sqrt{3}Le^{(12\phi_B - 1)^2}}{4\alpha}.\tag{28}$$

We can now use these two expressions to convert our computational/analytical findings to the experimentally available results for any value of solid volume fraction.

In their Section 2, Ziehmer et al. [33] report a volume fraction of $\phi_B = 0.296$ for as-prepared nanoporous gold. In Ref. [24], they determine a scaled connectivity density of $g^2 = 0.124$. This can be converted to the current notation, say $g^{M_1/S}$, using

$$g^{M_1/S} = g^2 \times \left[ \frac{L}{T} \right]^3.\tag{29}$$

In view of Eq. (27) with $\alpha \approx 1.23$, for $\phi_B = 0.296$, one has $L = 0.461L$. Thus, with our scaling, $g^{M_1/S} \approx 0.124/0.461 = 1.267$. For $\phi_B = 0.296$, the prediction of the scaled genus density, making use of Eq. (16), is $g_x^2 \approx 1.389$. Thus, we have $g^{M_1/S}/g_x^2 \approx 0.91$.

In their Table 1, Mangipudi et al. [23] report volume fractions of $\phi_{B,1} = 0.30 \pm 0.1$ and $\phi_{B,2} = 0.32 \pm 0.5$ for their nanoporous gold samples, for which the scaled connectivity densities are $g^{M_1} = 0.067$ and $g^{M_2} = 0.043$, respectively. These can be converted to the current notation, say $g^{M_{i/S}}$ for $i = 1, 2$, using

$$g^{M_{i/S}} = g^{M_i} \times \left[ \frac{L}{S_{i-1}} \right]^3.\tag{30}$$

In view of Eq. (28) with $\alpha \approx 1.23$, for $\phi_{B,1} = 0.30$, one has $L/S_1 \approx 2.476$. Thus, with our scaling, $g^{M_{1/S}} \approx 0.067 \times 2.476 = 1.017$. For $\phi_B = 0.30$, the prediction of the scaled genus density, making use of Eq. (16), is $g_x^2 \approx 1.422$. For $\phi_{B,2} = 0.32$, one has $L/S_2 \approx 2.546$. Thus, with our scaling, $g^{M_{2/S}} \approx 0.043 \times 2.546 = 0.710$. The prediction of the scaled genus density, making use of Eq. (16), is $g_x^2 \approx 1.576$. Thus, we have $g^{M_{1/S}}/g_x^2 \approx 0.71$ and $g^{M_{2/S}}/g_x^2 \approx 0.45$.

Appendix B. Supplementary data

Supplementary data related to this article can be found at https://doi.org/10.1016/j.actamat.2018.01.005.

References


