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Thermoelectric coefficients of n-doped silicon from first-principles via the solution of the Boltzmann transport equation

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We present a first-principles computational approach to calculate thermoelectric transport coefficients via the exact solution of the linearised Boltzmann transport equation, also including the effect of non-equilibrium phonon populations induced by a temperature gradient. We use density functional theory and density functional perturbation theory for an accurate description of the electronic and vibrational properties of a system, including electron-phonon interactions; carriers’ scattering rates are computed using standard perturbation theory. We exploit Wannier interpolation (both for electronic bands and electron-phonon matrix elements) for an efficient sampling of the Brillouin zone, and the solution of the Boltzmann equation is achieved via a fast and stable conjugate gradient scheme. We discuss the application of this approach to n-doped silicon. In particular, we discuss a number of thermoelectric properties such as thermal and electrical conductivities of electrons, Lorenz number and Seebeck coefficient, including the phonon drag effect, in a range of temperatures and carrier concentrations. This approach gives results in good agreement with experimental data and provides a detailed characterization of the nature and the relative importance of the individual scattering mechanisms. Moreover, the access to the exact solution of the Boltzmann equation for a realistic system provides a direct way to assess the accuracy of different flavours of relaxation time approximation, as well as of models that are popular in the thermoelectric community to estimate transport coefficients.

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I. INTRODUCTION

A detailed understanding of electrical transport and energy dissipation phenomena is crucial for the discovery and the development of high-performance materials and devices for applications ranging from nanoelectronics to energy conversion technologies.

For this, there is nowadays a growing demand for accurate and efficient computational tools to compute electrical transport coefficients from first-principles. The challenge here is a proper description of the dynamics of carriers, while accounting for the relevant carriers’ scattering mechanisms. In this context, the Boltzmann transport equation (BE) offers a convenient framework for a detailed microscopic description of transport in metals and semiconductors. However, achieving the full predictive power of this theory requires the exact solution of the BE, that is a complex integro-differential equation, as well as accurate materials’ parameters, including electronic band structures and electron-phonon and electron-defect scattering terms.

Different flavours of relaxation-time approximation are often used to tackle the complexity of the BE, but their accuracy in predicting the changes of properties with temperature or carrier concentration is not always satisfactory: for instance, Xu and Verstraete have recently shown that the constant relaxation-time approximation, a popular and quite successful scheme to estimate electronic transport coefficients, provides qualitatively wrong results for the Seebeck coefficient of a relatively simple system such as bulk lithium.

Only in recent years the first efforts have been put forward to solve numerically the BE beyond the relaxation-time approximation and using ab initio materials’ parameters. For instance, Wang and coworkers computed the thermoelectric properties of n-doped silicon from the numerical solution of the linearized Boltzmann equation where the electronic band structure and average intervalley deformation potentials were computed from first-principles, while semiempirical models were used for intravalley scattering, charged impurity scattering and electron-plasmon coupling. In a more recent work, Li studied the transport properties of n-doped Si, Al and n-doped MoS by solving the linearised BE using a standard iterative technique and including ab initio bands and electron-phonon scattering obtained from a linear interpolation of the ab initio electron-phonon matrix elements computed on a coarse grid. Despite these advances, the calculation of thermo-electric transport properties fully from-first-principles still represents a challenge, even in the case of simple metals or semiconductors.

In this paper we present a first-principles computational infrastructure to calculate electrical transport coefficients of materials from the exact solution of the linearized BE. We use density functional theory (DFT) and density functional perturbation theory (DFPT) for an accurate description of the electronic and vibrational properties of a system, including electron-phonon interactions; carriers’ scattering rates are computed using standard perturbation theory. We exploit Wannier interpolation for both electronic bands and electron-phonon matrix elements for an efficient sampling of the Brillouin zone, and the solution of the Boltzmann equation is achieved via a fast and stable conjugate gradient scheme.

We discuss the application of this general methodology to n-doped silicon. In particular, we discuss a number of thermo-electric properties such as thermal and electrical conductivities of electrons, Seebeck coefficient and Lorenz number in a range of temperatures and donor concentrations, also including the effects related to nonequilibrium phonon populations induced by a temperature gradient.

Silicon represents an excellent test case to validate our ap-
proach as its electronic transport properties have been extensively studied in the past. In particular, there are several theoretical works focussing on transport coefficients of silicon computed using semiempirical models (see, for instance, Ref. 7) as well as using ab initio scattering terms within the relaxation time approximation.\textsuperscript{8–11} or within the exact solution of the BE.\textsuperscript{3,4} In addition to this, the role of the coupled electron-phonon dynamics on the thermoelectric properties of silicon has very recently attracted renewed interest.\textsuperscript{11–13} Indeed, in this material there is a substantial enhancement, even at room temperature and in heavily doped samples, of the Seebeck coefficient induced by the drag exerted on charge carriers from phonons diffusing along a temperature gradient (phonon drag). In the past this topic has not received much attention beyond the initial experimental\textsuperscript{14} and theoretical\textsuperscript{15} works, but with current interest in thermoelectric energy conversion this effect might represent an interesting route to enhance thermoelectric performance.\textsuperscript{11}

The structure of the paper is as follows. In the Sec. II, we present the theoretical framework for transport. The results for the transport coefficients on n-doped silicon are presented and discussed in Sec. III. The details of the BE solver, the convergence and the performance of the numerical techniques are discussed in the appendices.

II. METHODOLOGY

A. Boltzmann transport in linear response

When an electric field $\mathbf{E}$ and a temperature gradient $\nabla_r T$ are present across a material, the steady-state distribution function $f_m(k)$ of charge carriers on band $m$ and wave vector $k$ can be described by the BE\textsuperscript{16}

$$v_{mk} \cdot \nabla_r T \frac{\partial f_m(k)}{\partial T} - \frac{e}{h} \mathbf{E} \cdot \nabla k f_m(k) = \left( \frac{\delta f_m(k)}{\delta t} \right)_{\left(\text{coll}\right)},$$

(1)

where $e$ is the electronic charge, $v_{mk} = \sqrt{\frac{k_B T}{\mu}}$ is the electron velocity (where $\mu$ is the quasiparticle energy). The term on the right is the collision term given by

$$\left( \frac{\delta f_m(k)}{\delta t} \right)_{\left(\text{coll}\right)} = \frac{1}{N_k} \sum_{m'k'} \left\{ f_{m'}(k') W_{m'm,k} - f_m(k) W_{m,k'm'} \left[ 1 - f_m(k') \right] \right\},$$

where $N_k$ is the number of k-points in the Brillouin zone (BZ) and $W_{m,k'm'}$ is the probability per unit time for an electron in a state $m, k$ to be scattered into an empty state $m', k'$ (the details of $W_{m,k'm'}$ for the different scattering mechanisms are given in Sec. II.B)

If the external fields are weak, it is possible to write the steady-state distribution in terms of the equilibrium one, $f_m^0(k)$, and a term that is linear in the perturbation fields,

$$f_m(k) = f_m^0(k) - \sum_{\epsilon_m} \frac{\partial f_m^0(k)}{\partial \epsilon_m} \chi_m(k) (2)$$

In this way it is possible to derive the linearised Boltzmann equation (BE) that can be solved to obtain $\chi_m(k)$

$$- \frac{\partial f_m^0(k)}{\partial \epsilon_m} v_{mk} \cdot \left[ E e + \nabla_r T \left( \frac{\epsilon_m - \mu}{T} \right) \right] = \frac{1}{k_B T N_k} \sum_{m'k'} \Pi_{m'k',mk}^0 \left[ \chi_m(k') - \chi_m(k) \right]$$

(3)

where $\mu$ is the chemical potential and we have used $\Pi_{m'k',mk}^0 = f_m^0(k') \left[ 1 - f_m^0(k) \right] W_{m'm,kk'}$. This equation can be written in the matrix form $b = Ax$, where $b$ is the left-end side of Eq. 3, $x$ is $\chi_m(k)$ and the matrix $A$ is

$$A_{ij} = \frac{1}{k_B T N_k} \left[ \Pi_{ij} - \delta_{ij} \sum_{m'} \Pi_{mm'}^0 \right]$$

where $i = \{m, k\}$, $A$ is symmetric and positive semidefinite.\textsuperscript{16}

The expressions for the charge current, $j_e$, and the heat current, $j_c$, are

$$j_e = \frac{2e}{VN_k} \sum_{mk} v_{mk} \frac{\partial f_m^0(k)}{\partial \epsilon_m} \chi_m(k),$$

(4a)

$$j_c = -\frac{2}{VN_k} \sum_{mk} \left( \epsilon_m - \mu \right) v_{mk} \frac{\partial f_m^0(k)}{\partial \epsilon_m} \chi_m(k),$$

(4b)

where $V$ is the volume of the unit cell. In linear response the electric and thermal current densities are related to the external fields via transport coefficients

$$j_e = \sigma \left[ E + \nabla_r T \left( \frac{-\nabla_r T}{T} \right) \right],$$

(5a)

$$j_c = \alpha E + \beta T \left( \frac{-\nabla_r T}{T} \right).$$

(5b)

It has to be remarked that in this formalism $E = \nabla_r (\phi - \nu / \epsilon)$ is the effective electric field arising from both an electric and chemical potential gradient and should be regarded as the electrochemical field. Also, in principle the transport coefficients are tensors,\textsuperscript{17} but they are scalars in the case of silicon.

The quantities on which we will focus on for n-doped silicon are: the electrical resistivity $\rho_e = 1/\sigma$ (where $\sigma$ is the electrical conductivity), the electron mobility $\mu_e = e\sigma / n$ (where $n$ is the carrier concentration), the Seebeck coefficient $S$ and the Lorenz number $L = \hbar / eT$ (where $\hbar = \beta - \alpha S$ is the thermal conductivity of the electrons).

In order to establish a link with the relaxation-time approximation, it is possible to rewrite $\chi_m(k)$ in this way:

$$\chi_m(k) = -v_{mk} \cdot \left[ E e + \nabla_r T \left( \frac{\epsilon_m - \mu}{T} \right) \right]$$

(6)

where we have introduced two effective transport relaxation times $\tau_m^e(k)$ and $\tau_m^i(k)$ that are determined from the solution
of the BE’s that include only the electric field and only the thermal gradient, respectively. With this definition, the transport coefficients have familiar forms:

\[ \sigma = \frac{2e^2}{VN} \sum_{mk} \frac{\partial f^0_m(k)}{\partial \epsilon_m} [v_m^\parallel(k)]^2 e_m^\parallel(k) \]  

(7a)

\[ \sigma S = \frac{2e}{TVN} \sum_{mk} \frac{\partial f^0_m(k)}{\partial \epsilon_m} [v_m^\parallel(k)]^2 (\epsilon_m - \epsilon_t) e_m^\parallel(k) \]  

(7b)

\[ a = \frac{2e}{VN} \sum_{mk} \frac{\partial f^0_m(k)}{\partial \epsilon_m} [v_m^\parallel(k)]^2 (\epsilon_m - \epsilon_t) e_m^\parallel(k) \]  

(7c)

\[ \beta = -\frac{2}{TVN} \sum_{mk} \frac{\partial f^0_m(k)}{\partial \epsilon_m} [v_m^\parallel(k)]^2 (\epsilon_m - \epsilon_t)^2 e_m^\parallel(k) \]  

(7d)

where \( v_m^\parallel(k) \) is the component of the velocity along the applied field.

### B. Sources of electronic scattering

The main ingredients that determine the transport characteristics of a material are the transition probabilities per unit time, \( W_{mk \rightarrow m'k+q} \) on the right-hand side of Eq. 3. Below we discuss the transition probabilities for the sources of electronic scattering that we used to simulate the transport properties of n-doped silicon.

#### 1. Electron-phonon coupling

Under the assumption that lattice vibrations are in thermal equilibrium, the transition probability due to electron-phonon scattering is given by:

\[ W_{el-phon}^{el-imp} = \frac{2\pi}{\hbar} \sum_{\lambda} \left| g^\lambda_{mn}(k; q) \right|^2 \]

\[ \times \left\{ n_{\lambda}^0(q) \delta[\epsilon_{m'k+q} - \epsilon_m - \hbar \omega_{\lambda}(q)] 
+ \left[ n_{\lambda}^0(q) + 1 \right] \delta[\epsilon_{m'k+q} - \epsilon_m + \hbar \omega_{\lambda}(q)] \right\} \]  

(8)

where \( \omega_{\lambda}(q) \) and \( n_{\lambda}^0(q) \) are the phonon frequency and the equilibrium distribution of the phonon mode of branch index \( \lambda \) and wave vector \( q \), and the electron-phonon matrix elements are given by:

\[ g_{mn}(k; q) = \sqrt{\frac{\hbar}{2M\omega_{\lambda}(q)}} (mk + q) \partial_{\lambda} V |m'k] \]  

(9)

where \( |m'k] \) and \( |mk + q] \) are the Bloch eigenstates, \( M \) is the atomic mass and \( \partial_{\lambda} V \) is the derivative of the Kohn-Sham self-consistent potential with respect to the atomic displacement corresponding to a phonon mode \( \lambda \) at \( q \).

#### 2. Ionized impurities

Even though it is nowadays possible to compute form first-principles the interaction between charge carriers and ionized impurities in semiconductors, in this work we rely on a simple analytical model developed by Brooks and Herring (BH). In this model the transition probability per unit time is given by:

\[ W_{el-imp}^{el-imp} = \frac{2\pi Z^2 n_i e^4}{\hbar^2 (\epsilon_{m'k+q}^2 - \epsilon_m^2)} \left( \beta_z^2 + |q|^2 \right)^2, \]

(10)

where \( Z \) is the charge of the impurity, \( n_i \) is the impurity density, \( \epsilon_t \) and \( \epsilon_{0} \) are the relative and the vacuum permittivity. In the Debye approach, for non-degenerate semiconductors, the reciprocal of the screening length, \( \beta_z \), is given by:

\[ \beta_z = \frac{\epsilon_{0}}{\epsilon_{0} q_{0} k_B T}. \]

(11)

In this work we focus on n-doped silicon and assume that all the donor impurities are completely ionised, i.e. the donor density is equal to the free carriers’ concentration.

### C. Phonon drag

The assumption of thermal equilibrium of the lattice is not strictly valid anymore when a temperature gradient is present across a material. In this case, non-equilibrium phonon populations can leave a fingerprint on the thermoelectric coefficients. This effect was first predicted by Gurevich in 1945, and only later, experiments on germanium and silicon clearly revealed the related anomalous peak in the Seebeck coefficient at low temperatures.

In order to accurately account for this effect, in principle one should solve the coupled BE’s for the electrons and phonons. An alternative approach to this computationally demanding effort was proposed by Cantrell et al. in the context of a relaxation time approximation and it exploits the possibility to partially decouple the electron and phonon transport for weak electron-phonon coupling. This approach has been recently implemented and used in silicon within an ab initio framework. In this way we proceed along a similar line but in the context of the exact solution of the BE. In particular, we partially decouple the two BE equations under the assumption that the non-equilibrium phonon populations that arise in presence of a temperature gradient can be computed neglecting the electron-phonon scattering in the BE for phonons. This approximation is reasonable in silicon at carriers’ concentration lower than \( 10^{19} \) cm\(^{-3} \) as in this regime the thermal conductivity is the same as in undoped silicon. In this way it is possible to compute the steady-state correction to the equilibrium phonon populations, \( \delta n_{\lambda}(q) \), from the BE for phonons, and include them into the full BE for the electrons. By considering only the first order deviation of the distribution functions, it is possible to obtain a linearised BE for electrons that is the...
same as Eq. 3 apart an additional term on the left-hand side given by

\[
\left( \frac{\delta f_m(k)}{\delta t} \right)_{\text{Ph drag}} = -\frac{2\pi}{\hbar} \sum_{q,m'} |g_{mn}^k(k;q)|^2 \\
\times \left\{ \delta n_k(q) \delta[\epsilon_{m'k+q} - \epsilon_{mk} - \hbar\omega_{k}(q)] \\
+ \delta n_k(-q) \delta[\epsilon_{m'k+q} - \epsilon_{mk} + \hbar\omega_{k}(q)] \right\} [f^0_{m}(k + q) - f^0_{m}(k)]
\]

This term is the driving force that is responsible for an additional electrical current in the same direction as the heat flow, as if the charge carriers were dragged along by phonons. This effect results in a phonon drag contribution to the Seebeck effect that adds to the standard diffusive part. In this work we compute \( \delta n_k(q) \) using the single-mode relaxation time approximation

\[
\delta n_k(q) = -\tau_q(q) \frac{c_q(q) \cdot \nabla n_k^0(q)}{k_BT^2} \hbar \omega_k(q) [n_k^0(q) + 1]
\]

where \( \tau_q(q) \), \( c_q(q) \) are the anharmonic phonon lifetime and the group velocity of the phonon mode \( \lambda \) at \( q \), respectively.

D. Computational details

We used DFT and DFPT as implemented in the Quantum-ESPRESSO distribution\(^{26}\) within the local-density approximation (LDA)\(^{27}\) to compute the electronic and vibrational properties including the electron-phonon coupling matrix elements. We used a norm-conserving pseudopotential and a plane-wave expansion up to a 30 Ry cutoff. Brillouin-zone sampling was performed on 12×12×12 Monkhorst-Pack mesh for all charge density and phonon calculations. We used a theoretical lattice parameter of 5.41 Å. All the quantities \( \epsilon_{mk} \), \( v_{mk} \), and \( g_{mn}^k(k; q) \) have been first calculated on grids of 10×10×10 k-points and 5×5×5 q-points. Then we have used a first-principles interpolation scheme\(^{26}\) based on maximally localized Wannier functions,\(^{5}\) as implemented in the Wannier90\(^{28}\) and EPW\(^{29}\) packages to compute transport quantities on very dense meshes up to 110×110×110. Wannier interpolation has been used successfully in recent works on the electrical transport properties of graphene,\(^{30}\) phosphorene,\(^{31}\) silicon\(^{11}\) and bulk metals.\(^{32}\)

Even though the Wannier interpolation of the electron-phonon matrix elements is computationally efficient, it is quite time consuming when many k-points are required to achieve converged transport coefficients. To reduce the computational burden, we have adopted a double grid scheme, in which the electron-phonon coupling matrix elements computed via Wannier interpolation are stored on a fine mesh in reciprocal space (\( N_{EPC} \times N_{EPC} \times N_{EPC} \)), and an ultra fine grid (\( N_f \times N_f \times N_f \), with \( N_f = p \cdot N_{EPC} \) and \( p = 1, 2, 3 \)) is used to solve the BE and compute the transport quantities. For the points in the ultra fine grid we only compute electronic bands and phonon frequencies, and use the electron-phonon matrix elements at the closest points in the \( N_{EPC} \times N_{EPC} \times N_{EPC} \) grid. This choice is motivated by the fact that the electron-phonon matrix elements are a rather smooth function in the Brillouin zone while the transition probabilities (involving delta functions and products of distribution functions) exhibit a more oscillatory behavior. The delta functions in the transition probabilities are replaced with gaussians. Converged transport properties were obtained with \( N_{EPC} = 70 \), \( p = 3 \) and a gaussian width of 4 meV. Doping has been simulated within the rigid band approximation via a shift of the chemical potential. The convergence of the transport properties with respect to the grids and gaussian width is discussed in App. B. The solution of the BE is achieved with an in-house parallel code that uses a fast and stable preconditioned Conjugate Gradient (CG) algorithm.\(^{33}\) A similar approach has been used for the solution of the BE for phonons by Fugallo et al.\(^{34}\). App. A discusses the accuracy and the performance of the BE solver.

The anharmonic phonon lifetimes needed to compute the contribution of the phonon drag to the transport coefficients were computed within the single mode relaxation time approximation and from the third order force constants used in Refs. 35 and 36. In this work we have recalculated the lifetimes on a \((70 \times 70 \times 70)\) mesh of q-points using a gaussian width of 0.25 meV.

III. RESULTS

A. Mobility and resistivity

In Figs. 1 and 2 we present our results for the electron mobility of n-doped silicon. The room temperature results (Fig. 1) show that for carrier concentrations lower than 10\(^{15}\) cm\(^{-3}\), the mobility is mainly limited by electron-phonon interactions; at higher carriers concentrations, instead, the mobility decreases as a result of an increased electron-impurity scattering. As the temperature increases the mobility decreases (Fig. 2) mainly because of a stronger electron-phonon scattering. Our predicted electron mobility is in reasonable agreement with the reported experimental data as a function of doping and temperature. It is however clear that our results tend to overestimate the experimental results (but note that the uncertainty in the experimental data is quite large, around 10% at low doping). This could be related to an underestimation of the electron-phonon coupling computed within standard DFT at the level of LDA. Indeed, similar shortcomings of LDA have been reported quite recently for other materials, including simple sp-bonded compounds, such as graphene\(^{30}\) and diamond.\(^{38}\) It is important to point out that the discrepancy at high doping with the experimental data could also be related to additional scattering mechanisms involving electron-plasmon interactions\(^{39}\) that we have not included in this work.

It is interesting to compare our results with previous DFT-based theoretical work. Focussing on the room temperature, low doping regime our prediction for the phonon-limited mobility is 1750 cm\(^{2}\)V\(^{-1}\)s\(^{-1}\). Our result is slightly closer to experiments than the value of 1860 cm\(^{2}\)V\(^{-1}\)s\(^{-1}\) reported in Ref. 4, where the BE is numerically solved. This difference is
probably related to the different interpolation schemes (Wannier vs linear interpolation) of the electron-phonon matrix elements and the result of a different sampling of the BZ. The comparison with DFT-based results obtained within the relaxation time approximation shows that our result is somehow in between the reported values of 1970 cm$^2$V$^{-1}$s$^{-1}$ (see Ref. 8) and around 1550 cm$^2$V$^{-1}$s$^{-1}$ (see Ref. 11). While the discrepancies with respect to Ref. 8 could be related to a different sampling of the BZ, it is not easy to identify the origin of the differences with the results of Ref. 11, in which both the interpolation scheme and the k-point grids used are very close to the ones used in this work. However, it is interesting to point out that, as shown in App. C, the use of the relaxation time approximation as in Ref. 11 gives electron mobilities that are 4-5% lower than the prediction from the exact solution of the BE, making the discrepancy between the two works slightly less pronounced.

It is interesting at this point to analyse in more detail the contribution of the different scattering channels to the electrical transport. As shown in Fig. 1, information about the role of different mechanisms can be obtained switching on and off interactions when solving the BE. This analysis however is only qualitative as the effects are not additive. In order to provide a quantitative estimate of the contributions from the different scattering channels, a possibility is to focus on the electrical resistivity and exploit the variational formula

$$\rho_e = \frac{V}{4e^2k_BT} \sum_{m,k,m',k'} \left[ \chi_m(k) - \chi_m(k') \right]^2 \Pi_{m,k,m',k'}^0 \left[ \sum_{m,k} V_{m,k}(k) \chi_m(k) \frac{\partial \Pi_{m,k}(k)}{\partial \chi_m(k)} \right]^2$$

This formula has its minimum value (the steady-state electrical resistivity) when $\chi_m(k)$ is the solution of the BE, and it is often used to determine an upper bound to the resistivity when using trial functions. In this case, however, we insert in Eq. 14 the solution of the BE and exploit the fact that in Eq. 14 the different contributions from the scattering mechanisms are additive. The electrical resistivity at room temperature is shown in Fig. 3: our results reproduce well the experimental data as a function of carrier concentration. In Fig. 4 we present the results of the analysis based on Eq. 14 of the different contributions to the resistivity in the low and high doping regimes. At low doping, the role of impurities is negligible and the main contribution (around 55%) comes from intra-valley acoustic phonon scattering. The remaining 45% is from inter-valley transitions. It is interesting to observe that around 35% of these contributions are due to scattering with inter-valley optical phonon modes. At higher doping (in a regime in which screening effects are weak), charged impurity scattering becomes dominant, accounting for almost 70% of the resistivity. Inter-valley channels also increase their importance relatively to intra-valley ones.
The phonon drag adds up a nearly constant term that con-
mines the overall dependence on the carrier concentration.
The diffusive part of the phonon drag have been calculated using the single-mode
approximation implemented in codes such as Bolztrap44 and
Boltzmann.45 The reason for this is that the diffusive part of $S$
depends very little on the nature of the scattering mechanisms.
Indeed, as shown in Fig. 5, switching off phonon or impurity
scattering when solving the BE affects the result for $S$ by less
than 8%.

In order to analyse the contribution of the different phonon
modes to the phonon drag part of $S$, we have solved the BE
restricting the term in Eq. 12 to different phonon frequency
windows and to specific phonon modes. The results of this
analysis are shown in Fig. 6. Our results indicate that low-
frequency acoustic phonon modes play a crucial role. For in-
stance, Fig. 6 shows that phonons with frequency below 50
$\text{cm}^{-1}$ contribute to about 80% of the phonon drag effect, with
the longitudinal acoustic (LA) modes contributing almost 60% of the total. These data confirm previous findings obtained

FIG. 4. Relative contribution of the different scattering mechanisms
to electrical resistivity at $T = 300$ K. Carriers’ concentrations are
$2.4 \times 10^{14}\text{cm}^{-3}$ (upper panel) and $1.3 \times 10^{15}\text{cm}^{-3}$ (lower panel). The
charged impurities’ contribution (Imp), intra-valley phonon scatter-
ing (Intra) and inter-valley phonon scattering (f and g stand for f-
processes and g-processes, described in Ref. 20) are shown. The
phonon scattering are separated in acoustic (dark red) and optical
(light yellow) branches, and different patterns are used for the modes:
diagonal lines for TA1 and TO1, dense diagonal lines for TA2 and
TO2 and vertical lines for LA and LO. The phonon energies of the
different modes are also shown.

B. Seebeck coefficient

Fig. 5 shows the theoretical results for the Seebeck coeffi-
cient as a function of carrier concentration at $T = 300K$.
The diffusive part of $S$ is the largest contribution and deter-
mines the overall dependence on the carrier concentration.
The phonon drag adds up a nearly constant term that con-
tributes around 25% of the total $S$ at $T = 300K$.

At low doping our results slightly underestimate the mea-
sured $S$ but the overall agreement between theory and exper-
iment is good. The underestimation could be due to the fact that i) the out-of-equilibrium phonon populations that produce
the phonon drag have been calculated using the single-mode
relaxation time approximation and anharmonic force constants
computed on a $3\times3\times3$ supercell that tend to underestimate the
lattice thermal conductivity,36 and ii) as already mentioned,
our calculations seem to slightly underestimate the strength of the
electron-phonon coupling.

At very high doping, for $n \gtrsim 10^{19}\text{cm}^{-3}$, our results are
higher than the experimental data. This is due to the fact that we do not include the additional contribution of electron-
phonon scattering when computing the out-of-equilibrium phonon populations. As noted in Ref. 11, at high doping this
saturation effect15 of charge carriers on phonons results in a
substantial decrease of the phonon drag contribution and the
Seebeck coefficient approaches the diffusive part of $S$.

It is interesting to note that the diffusive part of $S$ agrees
quite well with the prediction of the constant relaxation time
approximation (CRTA), a popular and often quite successful
approximation implemented in codes such as Bolztrap44 and
Boltzwann.45 The reason for this is that the diffusive part of $S$
depends very little on the nature of the scattering mechanisms.
Indeed, as shown in Fig. 5, switching off phonon or impurity
scattering when solving the BE affects the result for $S$ by less
than 8%.
on the basis of the relaxation time approximation.\textsuperscript{11} As suggested in Ref. 11, an interesting strategy to enhance the $zT$ could be based on a phonon frequency filtering approach. For instance, in silicon at room temperature phonons with frequencies above 50 \text{ cm}^{-1} contribute to about 80\% of the thermal conductivity.\textsuperscript{11,35} Engineered scattering of phonons above 50 \text{ cm}^{-1} could therefore provide an effective way to strongly reduce the thermal conductivity while preserving the beneficial effect of low-frequency phonons on the Seebeck coefficient. This filtering approach could be based on impurity scattering that is typically more effective for phonons at higher frequency (see, for instance, Ref. 35 where this effect has been studied in SiGe alloys). This strategy and, in particular, the possibility to use nanoclusters as impurities is analysed in more details in Ref. 11.

![Figure 6](image6.png)

**FIG. 6.** Contribution of phonon modes to the phonon drag part of the Seebeck coefficient at $T = 300K$ and for $n = 1.75 \times 10^{14} \text{ cm}^{-3}$. The solid red line is the contribution from all the modes; the blue long-dashed (green short-dashed) line is the contribution from TA (LA) phonon.

Fig. 7 shows the Seebeck coefficient as a function of temperature at low doping. It is clear that the phonon drag becomes more important at low temperature, resulting in a larger correction to the diffusive $S$. As Fig. 6 shows, the disagreement between theory and experiment becomes more pronounced at low temperature. This is the result of the fact that the small underestimation that we pointed out for the results at 300 K becomes more marked at lower temperature. In particular, as shown in Ref. 36, our results for the lattice thermal conductivity of silicon show slightly larger deviations at low temperatures. A more accurate estimate of the Seebeck coefficient at low temperature would clearly require a more accurate calculation of phonon transport in silicon, but this is beyond the scope of this work.

![Figure 7](image7.png)

**FIG. 7.** Seebeck coefficient as a function of temperature for $n = 1.75 \times 10^{14} \text{ cm}^{-3}$. The red open circles are the the total theoretical Seebeck coefficient while the dashed red line is the diffusive part of $S$. The black diamonds are the experimental data.\textsuperscript{14}

![Figure 8](image8.png)

**FIG. 8.** Electronic thermal conductivity as a function of carrier concentration at $T = 300 \text{ K}$. The red solid line is from the solution of the BE that includes all the scattering mechanisms and the phonon drag. The blue long-dashed line is the same as the red solid line but without the phonon drag.

![Figure 8](image8.png)

**FIG. 8.** Electronic thermal conductivity as a function of carrier concentration at $T = 300 \text{ K}$. The red solid line is from the solution of the BE that includes all the scattering mechanisms and the phonon drag. The blue long-dashed line is the same as the red solid line but without the phonon drag.

C. Lorenz number

The Lorenz number (or Lorenz function) $L = \frac{\kappa_e}{\sigma T}$ is an important quantity in thermoelectrics research, as it provides a way to separate the electronic thermal and lattice thermal conductivity. Good thermoelectric compounds display a high degree of electronic and structural complexity, and it is not obvious that the Wiedemann-Franz law, derived for simple metals obeying elastic scattering, remains valid also in these compounds. As the design and the optimization of thermoelectric materials involve modifications to the electrical conductivity or lattice thermal conductivity, a precise knowledge of this quantity is important to accurately characterise and compare different compounds and determine successful routes towards higher thermoelectric performance.

The electronic contribution to the thermal conductivity, one of the ingredients to compute $L$, is shown in Fig. 8. This quantity increases with the carrier concentration, but even at $n = 10^{29} \text{ cm}^{-3}$ the room temperature electronic thermal conductivity is only a small fraction (about 2\%) of the lattice ther-
FIG. 9. Lorenz number $L$ as a function of carrier concentration at $T = 300$ K. The red solid line is from the solution of the BE that includes all the scattering mechanisms and the phonon drag. The red long-dashed line is the same as the red solid line but without the phonon drag. The green dot-dash-dashed (blue dot-dot-dashed) line is from the BE including only phonon (impurity) scattering. The orange short-dashed line is the CRTA. The black dotted horizontal line is the Wiedemann-Franz result.

FIG. 10. Lorenz number $L$ as a function of Seebeck coefficient $S$ at $T = 300$ K. The red solid line (dashed blue line) is from the solution of the BE with (without) the phonon drag. The green dot-dash-dashed line is the model proposed in Ref. 48. The black dotted horizontal line is the Wiedemann-Franz result.

The access to the exact solution of the BE for a realistic system allows the assessment of simplified models to extract the Lorenz number from experimental data. For instance, a recent work [48] proposes the use of an equation for $L$ in terms of the experimentally determined $S$: $L = 1.5 + \exp \left( \frac{S}{116} \right)$ (where $L$ is in $10^{-8}$ $W\Omega K^{-2}$ and $S$ in $\mu V/K$). This equation closely reproduces the $L$ obtained with a model that uses a single band with parabolic dispersion and an energy-dependent relaxation time accounting for acoustic phonon scattering. Fig. 10 shows the comparison of this model with our results. At low $S$ up to about 0.6 $mV K^{-1}$, a range for $S$ that is typical of good thermoelectrics, the model is not too far from the exact result and it certainly provides a better estimate to $L$ than the Wiedemann-Franz law. At higher $S$ the discrepancy becomes more pronounced as our predicted $L$ significantly decreases and reaches an asymptotic value at high-Seebeck of around $10^{-8}$ $W\Omega K^{-2}$.

IV. CONCLUSION

We have presented a detailed analysis of the thermoelectric coefficients of n-doped silicon via the exact solution of the linearised Boltzmann transport equation for electrons, also accounting for the effect of non-equilibrium phonon populations induced by a temperature gradient. We have discussed electron mobility, Lorenz number and Seebeck coefficient in a range of temperatures and carrier concentrations. The theoretical results compare reasonably well with the available experimental data and provide a detailed characterization of the relative importance of the different scattering mechanisms. We have also analysed the accuracy of different flavours of relaxation time approximations, as well as of a simplified model to extract the Lorenz number from experimental data. The computational approach that we have implemented is very general and can be applied to doped semiconductors as well to metallic systems. We believe that along this route more compounds will be systematically studied to gain a detailed microscopic understanding of their thermoelectric transport properties.
In order to solve the linear equation $Ax = b$ we use the preconditioned Conjugate Gradient method,\(^{33}\) where the diagonal of $A$ is used as preconditioner. The performance of the CG algorithm is shown in Fig. 11: the CG converges in 6 or 7 iterations, while with the addition of the preconditioning an accurate solution is obtained already after 2-3 iterations. In our experience standard iterative approaches\(^{49}\) converge very slowly, especially in the presence of charged impurity scattering.

In order to assess the accuracy of the solution of the BE, it is interesting to numerically verify the validity of the Kelvin relation, $\sigma S = a/T$. The lower panel of Fig. 11 shows that, when phonons are in thermal equilibrium, the converged solutions of the BE’s with temperature gradient and electric field as driving forces satisfy the Kelvin relation; the CG algorithm with preconditioning produces the converged result extremely quickly. In passing, it is worth mentioning that, as shown in Ref. 50, the Kelvin relations are satisfied also when the departure from equilibrium of both electron and phonon distribution functions is taken into account. In this work however we are not in the position to be able to verify directly this result as we have neglected the effect of the electrons on the phonon transport, an appropriate assumption only at low doping.

The BE solver implemented in this work has been fully parallelised for use on distributed memory clusters so that the collision matrix $A$, that scales quadratically with the number of electronic states, can be decomposed among the processes. An efficient way to tile the matrix is the chessboard fashion, by defining an MPI cartesian topology in which each process is given a tile of the matrix. The main advantage of the cartesian topology is that loading the collision matrix can be efficiently distributed, reducing greatly time and memory consumption.

**Appendix B: Sampling of the Brillouin zone**

Here we discuss the choice of the numerical parameters used to compute transport coefficients. These parameters are i) the mesh $N_{\text{EPC}} \times N_{\text{EPC}} \times N_{\text{EPC}}$ of k-points on which we compute the electron-phonon coupling matrix elements via Wannier interpolation, ii) the fine grid $N_f \times N_f \times N_f$ (where we use $N_f = p \cdot N_{\text{EPC}}$, with $p = 1, 2$ and 3) used to solve the BE and compute the transport quantities, iii) the energy window from the bottom of the conduction band used to restrict the number of k-points, iv) the width $\sigma_g$ of the gaussians used to replace the delta functions in the transition probabilities.

In Fig. 12 we present the electron mobility as a function of the number of k-points used to solve the BE. The different electron-phonon coupling grids are used: $N_{\text{EPC}} = 70$ (red empty symbols), $N_{\text{EPC}} = 90$ (blue patterned symbols) and $N_{\text{EPC}} = 110$ (green full symbols). The values of $\sigma_g$ used are 1 meV (triangles), 4 meV (circles) and 7 meV (squares). The results obtained with less (more) than $10^4$ ($5 \times 10^4$) k-points on the x-axis are for $p = 1$ ($p = 3$); the remaining points are for $p = 2$. The energy window used is 0.15 eV.

"FIG. 12. Electron mobility at room temperature and for $n = 1.6 \times 10^{14}$ cm$^{-3}$ as function of number of k-points used to solve the BE. Different electron-phonon coupling grids are used: $N_{\text{EPC}} = 70$ (red empty symbols), $N_{\text{EPC}} = 90$ (blue patterned symbols) and $N_{\text{EPC}} = 110$ (green full symbols). The values of $\sigma_g$ used are 1 meV (triangles), 4 meV (circles) and 7 meV (squares). The results obtained with less (more) than $10^4$ ($5 \times 10^4$) k-points on the x-axis are for $p = 1$ ($p = 3$); the remaining points are for $p = 2$. The energy window used is 0.15 eV."

**Appendix A: Performance and accuracy of the BE solver**

In order to solve the linear equation $Ax = b$ we use the preconditioned Conjugate Gradient method,\(^{33}\) where the diagonal of $A$ is used as preconditioner. The performance of the CG algorithm is shown in Fig. 11: the CG converges in 6 or 7 iterations, while with the addition of the preconditioning an accurate solution is obtained already after 2-3 iterations. In our experience standard iterative approaches\(^{49}\) converge very slowly, especially in the presence of charged impurity scattering.

In order to assess the accuracy of the solution of the BE, it is interesting to numerically verify the validity of the Kelvin relation, $\sigma S = a/T$. The lower panel of Fig. 11 shows that, when phonons are in thermal equilibrium, the converged solutions of the BE’s with temperature gradient and electric field as driving forces satisfy the Kelvin relation; the CG algorithm with preconditioning produces the converged result extremely quickly. In passing, it is worth mentioning that, as shown in Ref. 50, the Kelvin relations are satisfied also when the departure from equilibrium of both electron and phonon distribution functions is taken into account. In this work however we are not in the position to be able to verify directly this result as we have neglected the effect of the electrons on the phonon transport, an appropriate assumption only at low doping.

The BE solver implemented in this work has been fully parallelised for use on distributed memory clusters so that the collision matrix $A$, that scales quadratically with the number of electronic states, can be decomposed among the processes. An efficient way to tile the matrix is the chessboard fashion, by defining an MPI cartesian topology in which each process is given a tile of the matrix. The main advantage of the cartesian topology is that loading the collision matrix can be efficiently distributed, reducing greatly time and memory consumption.

"FIG. 11. Performance of the CG algorithm. Upper panel: resistivity (normalised to the converged result $\rho_c$) as a function of the number of iterations. Lower panel: Deviation from the Kelvin relation as a function of the number of iterations. The dashed blue (solid red) line is the result of the CG (with preconditioning). The results are for $n = 1.75 \times 10^{14}$ cm$^{-3}$, T=300 K and with phonons in thermal equilibrium."

"Appendix A: Performance and accuracy of the BE solver"
FIG. 13 clearly shows that the accurate evaluation of the Seebeck coefficient requires an energy window considerably smaller than the one needed to compute the Lorenz number. In addition, the energy window might also depend slightly on the relevant scattering mechanisms. It is possible to see this for the Lorenz number (a quantity that is much more sensitive than the Seebeck coefficient to the scattering mechanisms included in the BE): when only impurity scattering is considered there is large contribution from electronic states quite far from the bottom of the conduction band and the energy window required for the convergence is slightly larger than in the case in which only phonon scattering is included.

Appendix C: Approximate approaches

Fig. 14 compares different approximate approaches within the BE framework to estimate the phonon-limited electron mobility at low doping. A popular approach is the relaxation time approximation (RTA) in which a relaxation time is computed using the formula

\[
\frac{1}{\tau_{RTA}^{mk}} = \frac{2\pi}{h} \sum_{q,m',\lambda,q} |g_{n\lambda m'}^{q}(k;q)|^2 
\times \left\{ [f_{0 \lambda}^{0}(k+q) + n_{\lambda}^{0}(q)] \delta \left[ \epsilon_{m'k+q} - \epsilon_{mk} - \hbar \omega_{\lambda}(q) \right] 
\right.
\]

\[
+ [1 + n_{\lambda}^{0}(q) - f_{0 \lambda}^{0}(k+q)] \delta \left[ \epsilon_{m'k+q} - \epsilon_{mk} + \hbar \omega_{\lambda}(q) \right] \right\} 
\]

(C1)

and it is then used in Eq. 7 to compute the transport coefficients. As shown in Fig. 14, this approach gives mobilities that are around 4-5 % lower than the exact solution of the BE. It is interesting to use \( \chi_{\sigma}(k) \propto v_{n}^{(n)}(k) \tau_{RTA}^{mk} \) as a trial function in the variational formula for the resistivity, Eq.14. Also in this case the mobility is lower than the exact result (because of the variational principle), but now the difference with respect to the exact result is less than 1%. A simpler approach, in which the trial function is only proportional to the velocity gives results that are fortuitously close to the experimental results but much lower than the result obtained from the exact solution of the BE.

FIG. 14. Intrinsic electron mobility as a function of temperature computed with different approaches: the solid red line is from solution of the BE (CG), the dashed green line is the relaxation time approximation (RTA), the dot-dashed blue line (dotted orange line) is obtained using \( \chi_{\sigma}(k) \propto v_{n}^{(n)}(k)\tau_{RTA}^{mk} \) as trial function in the variational formula Eq.14. The symbols are the experimental data as in Fig. 2.

FIG. 13. \( L \) and \( S \) as a function of the energy window used to compute the integrals of the transport coefficients in Eq. 7. Panel above: relative accumulation of \( S \) (solid blue line) and \( L \) (dashed red line) computed under the assumption of phonons in thermal equilibrium. Panel below: \( L \) as a function of the energy window for different scattering mechanisms. Dashed red line corresponds to \( L \) computed with all scattering terms, while the solid green line (dotted blue line) is for the impurity (phonon) contribution. In all cases the BE at room temperature and for \( n = 10^{14} \text{ cm}^{-3} \) is solved for all the states up to 0.32 eV above the bottom of the conduction band.

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This energy window includes also the contribution of the second conduction band that starts around 0.13 eV from the bottom of the conduction band.