Superconductivity in diamond, electron–phonon interaction and the zero-point renormalization of semiconducting gaps

Manuel Cardona*

Max-Planck-Institut fuer Festkoerperforschung, Heisenbergstrasse 1, 70569 Stuttgart, Germany

Received 26 January 2006; received in revised form 7 March 2006; accepted 13 March 2006

Available online 1 August 2006

Abstract

The superconductivity observed in boron-doped diamond is discussed from the point of view of the phonon-driven BCS theory. It is shown that electron–phonon interaction is particularly strong in diamond. Other semiconductors possibly exhibiting superconductivity are brought to the fore. Related evidence for strong electron–phonon interaction in semiconductors containing carbon, nitrogen and oxygen is presented.

Keywords: A. Semiconductors; A. Superconductivity in diamond; D. Electron–phonon interaction; D. Energy gaps

1. Introduction

During the workshop definitive evidence of the appearance of superconductivity in heavily doped p-type diamond has been presented. After the initial discovery by Ekimov et al. [1], using boron-doped bulk diamond samples prepared under high-pressure and -temperature, the phenomenon has been confirmed for chemical vapor deposition (CVD) films [2] prepared by (CVD). The high pressure and temperature technique requires rather specialized equipment. Hence, superconducting bulk diamond samples prepared by this technique have thus far only been prepared in Ekimov’s group. The CVD technique, however, is easily available and most recent work dealing with superconductivity in boron-doped diamond has been performed on samples prepared by this technique (see articles in these proceedings). Critical temperatures $T_c$ as high as 11.4 K have been reported at this workshop by Ishizaka et al., this $T_c$ being limited by the amount of substitutional (i.e., electrically active) boron present in the films.

The mechanism of the superconductivity in p-type diamond has also received considerable attention at the workshop. A number of calculations, both ab initio and semi-empirical, suggest that the mechanism is rather straightforward BCS pairing induced by the electron–phonon interaction. This interaction seems to be particularly large when involving valence electrons and optical phonons in materials containing atoms of the second row of the periodic table. Some publications expounding this mechanism had already appeared before the workshop [3–5]. A recent theoretical publication [6] suggests that the electron–phonon mechanism may also induce superconductivity in n-type diamond (e.g., phosphorus doped) with electron concentrations similar to those for which superconductivity has been found for holes in p-type material. Such electron concentrations do not seem to have been reached thus far but may be reachable in the future [7].

Although theorists strongly favor the simple phonon-induced BCS model for the superconductivity of p-type diamond, an alternative model has been discussed at the workshop by Baskaran [8]. This model postulates that the superconductivity is related to the insulator–metal Mott transition that appears in the impurity band with increasing boron concentration. The author draws support from soft X-ray absorption and emission spectroscopy data.

* Tel: +49 711 689 1710; fax: +49 711 689 1712.
E-mail address: m.cardona@fkf.mpg.de.

1468-6996/$ - see front matter © 2006 NIMS and Elsevier Ltd. All rights reserved.
doi:10.1016/j.stam.2006.03.009
which suggest the existence of a narrow hole band immersed in the usual filled valence band [9]. However, very recent work by some of the authors of [9], involving angular resolved photoemission, reveal a standard k-dependent diamond-like band structure with a Fermi energy which corresponds to the boron doping. In any case, the hole concentration present in the superconducting samples is well above that of the Mott transition \( (2 \times 10^{20} \text{cm}^{-3}) [10] \) and slightly above that at which the impurity band disappears and degenerate hole conduction starts \( (5 \times 10^{20} \text{cm}^{-3}) [11] \).

However, the “smoking gun” for the phonon-BCS mechanism would be the observation of an isotope effect with \( T_c \propto M_c^{-1/2} \) (carbon has two stable isotopes with \( M_c = 12 \) and 13). Although isotope effect measurements are being attempted at this point, they are rather difficult because of the difficulty in keeping the hole concentration constant, or simply well defined, while \( M_c \) is being changed.

Undoped diamonds with various isotopic abundances\(^1\) have been grown by the high-pressure high-temperature technique. Synthetic diamonds with an abundance of \(^{12}\text{C}\) of 99.9\%, higher than the natural one, have a thermal conductivity even higher than that of natural diamond, especially at ~90 K [12], a fact that makes them interesting as substrate for high performance electronic circuits. Isotopically modified diamonds have also been used for the investigation of a wide variety of electron–phonon and phonon–phonon interaction effects which, at low temperatures, become weaker for the larger isotope masses because of the smaller zero-point vibrational amplitudes.\(^2\) The dependence on isotopic mass of diamond of properties such as the lattice constant [14], the phonon frequencies and their line widths [15] as well as its hardness [16] have been studied. The effect of isotopic disorder on the phonons has also been investigated [15].

Relevant to the present work is the electron–phonon interaction in the low-temperature region which, as already mentioned, extends, in diamond, to well above room temperature. There has been for some time considerable theoretical [17] and some experimental evidence [18] suggesting that the deformation potentials representing the electron–phonon interaction are much larger in diamond than in other diamond-like semiconductors (e.g., Si, Ge). This evidence refers mainly to the interaction of optical phonons with electrons (or holes) at the top of the valence bands, the mechanism suspected to be responsible for the superconductivity in boron-doped diamond.

Another important consequence of the electron–phonon interaction is the temperature and isotope mass renormalization of the optical response functions, especially the imaginary part of the frequency-dependent dielectric function \( e_i (\omega) \) (i.e., the optical absorption spectrum). Such effects have been profusely investigated for a number of diamond, zincblende and wurtzite-type semiconductors, either by measuring the temperature dependence of structure in the \( e_i (\omega) \) spectra or, at low-temperature, by measuring the effects of varying the isotopic mass [13]. Since the sharpest structures involve transitions bringing an electron from the top of the valence band to an empty state at the bottom of the conduction band, the effects of electron–phonon interaction on them must bear an at least qualitative relationship with those responsible for the superconductivity in diamond. It has been conjectured that the large electron–phonon interaction is a property of semiconductors containing atoms of the second row of the periodic table (B,C,N,O), (i.e., materials such as diamond, boron nitride and, to a lesser extent, SiC, GaN and ZnO [5].

In this article, we review the deformation potentials representing the electron–phonon interaction at the top of the valence band and at the bottom of the conduction bands of diamond and compare them with those of silicon. We then discuss the details of the band structure of diamond (and silicon) relevant to BCS-type superconductivity and present a simple calculation of the critical temperature which falls on the same range as that calculated by more sophisticated ab initio techniques. The dependence of \( T_c \) on the hole concentration, as measured by Kačmarčík et al. [11] is compared to the prediction of our calculations and the possibility of superconductivity in heavily doped silicon is touched upon. We also present the experimental evidence for the unusually strong electron–phonon interaction acting on the electronic states responsible for the edge luminescence of diamond, GaN and ZnO.

Raman spectroscopy has been profusely used for the characterization of heavily doped diamond. Heavily boron-doped samples show a spectrum of the optical phonons (at \(~1330\text{cm}^{-1}\) ) with an asymmetry of the Fano type. This asymmetry results from intervalence band transitions, induced by the electron–phonon interaction, interfering with the phonon scattering amplitude, a phenomenon that has been studied in silicon in great detail [19]. The Raman phonon line shape in diamond contains information about the density of hole states as well as the electron phonon interaction [19–21] and conceivably could be used to extract information about \( T_c \). An adequate procedure has, however, not yet been found. It may be of interest to mention that the sign of the interference (i.e., constructive vs. destructive) observed for p-type diamond is the same as that of p-type silicon, i.e., it corresponds to a positive Fano parameter \( q \).

2. Valence band structure and hole–phonon interaction parameters

The top of the valence bands of diamond-like semiconductors occurs at \( k = 0 \) and is triply degenerate.\(^3\) The

\(^1\)The natural abundances are \(^{12}\text{C}: 99\%\), \(^{13}\text{C}: 1\%\).

\(^2\)Because of the high Debye temperature (~1800 K) the low-temperature region of diamond extends to above room temperature.

\(^3\)We neglect spin–orbit splitting which at \( \Gamma \) is 14meV for diamond, negligible compared with the Fermi energy \( E_F \) of superconducting samples.
three bands split when moving away from \( k = 0 \), the \( \Gamma \) point of the Brillouin zone. The so-called heavy hole bands remain doubly degenerate for \( k \) along \{100\} and \{111\} but they split along lower symmetry directions, including \{1\bar{1}0\}. A schematic diagram of the Fermi surface of diamond (also silicon) is shown in Fig. 1 in the (001) plane of \( k \)-space. Note the bulging out of this surface for the upper heavy hole band along the \{1\bar{1}0\} directions. This bulging out results from a large effective mass along \{1\bar{1}0\}. The three effective masses along \{1\bar{1}0\} can be calculated using the so-called Luttinger parameters \( \gamma_1, \gamma_2, \gamma_3 \) [22]:

\[
\begin{align*}
\gamma_1^{-1}(1\bar{1}0) &= \gamma_1 + \gamma_2 - 3\gamma_3, \\
\gamma_2^{-1}(1\bar{1}0) &= \gamma_1 - 2\gamma_2, \\
\gamma_3^{-1}(1\bar{1}0) &= \gamma_1 + \frac{2}{3} \gamma_2 + 2\gamma_3.
\end{align*}
\]

Replacing the values of \( \gamma_i \) given under “present work” in [22] into Eq. (1) we obtain \( m_1 = 0.23 \), \( m_2 = 0.36 \) and \( m_3 = 1.64 \). The large value of \( m_3 \) is responsible for the large bulging of the heavy hole Fermi surface seen in Fig. 1 along \{1\bar{1}0\}. Recent angular resolved photoemission measurements [23] reveal the three \( k \)-dependent valence bands with rather similar effective masses. The large mass \( m_3 \) of Eq. (1) does not appear because \{1\bar{1}0\} directions do not seem to have been investigated.

By averaging the masses along \{1\bar{1}0\}, \{111\} and \{100\} [22] one finds for the average density of states effective mass \( m^e_{\text{D}} = 1.9 \). With this mass and for a hole concentration \( N_h = 2 \times 10^{21} \text{cm}^{-3} \) we obtain, using the standard expressions for a parabolic band, \( E_F = 0.4 \text{eV} \), a Fermi energy which compares well with the value 0.6 eV given in [4] for \( N_h = 5 \times 10^{21} \text{cm}^{-3} \). The corresponding density of holes, which will be used in our estimates of \( T_c \), is 0.047 per eV, spin and primitive cell.

The hole–phonon interaction at the \( \Gamma \)-point is usually represented in the semiconductor literature [24] by a deformation potential \( d_0 \) (in eV), although sometimes the interaction constant \( D_1 = d_0/a_0 \) (in eV/cm, \( a_0 \) is the lattice constant) is used. \( d_0 \) is defined by keeping in mind that the valence bands at \( \Gamma \) split into a singlet and a doublet when a sublattice displacement \( u \) is applied along the [111] direction. The singlet–doublet splitting is written as

\[
\delta u_0 = \frac{3}{2} u d_0
\]

an expression which defines \( d_0 \). Values of \( d_0 \) between 50 and 100 eV are reported in [17] for diamond.

3. Estimated critical temperature \( T_c \)

For our “back-of-the-envelope” estimate we use the McMillan equation

\[
T_c = T_D \exp \left\{ - \left[ \lambda (1 + \lambda)^{-1} - \mu^* \right]^{-1} \right\},
\]

where \( T_D = 1942 \text{K} \) is the temperature which corresponds to the optical phonon frequency \( \omega_D = 1330 \text{cm}^{-1} \). Some authors use slightly lower values of \( T_D \) and \( \omega_D \) to include the softening or renormalization of the phonon frequencies induced by the electron–phonon interaction [3,19]. For the sake of simplicity, we use unrenormalized values of \( T_D \) and \( \omega_D \). In Eq. (3), the dimensionless electron–phonon interaction parameter is given by

\[
\lambda = N_d \langle D^2 \rangle / M \omega_D^2,
\]

where \( N_d \) is the density of holes at the Fermi surface (per Hartree, spin and unit cell), \( D \) the electron–phonon deformation potential (in Hartree/bohr), \( M \) the atomic mass in units of the electron mass. The units of energy in \( D \) and \( \omega_D \) must be chosen to be the same.

The bulging of the heavy hole surface along \{1\bar{1}0\}, as shown in Fig. 1, signifies that most holes occupy regions of \( k \)-space along the \{1\bar{1}0\} directions. The electron–phonon matrix elements \( D^2 \) should, in principle, be averaged around all the Fermi surfaces. Again for simplicity, we assume that they connect only heavy holes along \{1\bar{1}0\} (see thin lines in Fig. 1). The average of \( D^2 \) can then be performed by hand. We find for the average matrix element to be replaced into Eq. (3):

\[
\langle D^2 \rangle^{1/2} = (0.55 d_0/a_0) \text{eV/Å},
\]

We used for \( d_0 = 79 \text{eV} \) and \( \mu^* \approx 0 \) in Eqs. (3) and (4).\footnote{Some authors take \( \mu^* \approx 0.1 \) for the Coulomb repulsion parameter, but we prefer to take \( \mu^* \approx 0 \) because of the lack of detailed information.}

Replacing the values mentioned above we obtain \( \lambda \approx 0.2 \) and \( T_c = 4.2 \text{K} \) for a hole density \( N_h = 2 \times 10^{21} \text{cm}^{-3} \) (corresponding to a 2.3% concentration of substitutional B atoms). In Table 1 we compare this \( T_c \) with the values obtained by other authors.

Although limited to low \( N_B \) concentrations, the most comprehensive study of \( T_c \) vs. \( N_B \) has been reported in [27].
This study covers the $0 < T_c \leq 2.1$ K range. It is interesting to compare its results with those obtained from Eqs. (3) and (4) assuming all parameters constant except $N_d$ ($N_B$) which is taken to be that of a parabolic band approximation:

$$N_d = AN_h^{1/3},$$

where $A$ is an adjustable constant and $N_h$ is assumed to be equal to $N_B$. Fig. 2 shows a fit with Eq. (5) of the $T_c$’s reported in [27].

For $N_B$ lower than $5 \times 10^{20}$ cm$^{-3}$ the points calculated with Eq. (5) seem to be above the few experimental ones, thus signaling the approaching Mott transition. Above $N_B = 12 \times 10^{20}$ cm$^{-3}$ the only available experimental point lies below the calculated curve, thus suggesting the presence of electrically inactive interstitial boron atoms.

Superconductivity has been predicted for n-type diamond with values of $T_c$ vs. carrier concentration rather similar to those calculated for p-type samples [6]. Unfortunately, the solubility of the dopant (electron concentrations below $5 \times 10^{19}$ cm$^{-3}$ [7]) seems to be lower than that required for the appearance of superconductivity (3.6 $\times 10^{21}$ cm$^{-3}$ [6]). It is, however, easy to rationalize the similarity of the $T_c$’s predicted for p- and n-type diamond. In the n-type material the electrons occupy six valleys along the \{1 1 0\} directions, close to the X-points of the BZ. The longitudinal ($m_\parallel$) $m$ and transverse $m_\perp$ masses of these valleys are $m = 1.50$, $m_\perp = 0.34$ [22]. Taking into account the valley degeneracy, we find a density of states mass

$$\langle m_e^{3/2} \rangle = 6(m m_\perp^2)^{1/2} = 2.51$$

which is somewhat larger than that calculated for holes ($m_\perp^{1/2}$) = 1.91 leading, everything else being equal, to a larger $T_c$. The average electron–phonon deformation potential to be inserted in Eq. (4) is $8 \times 10^8$ eV/cm, somewhat lower than that used for p-type diamond ($12 \times 10^8$ eV). Hence, the product of $N_d$ times $D$ is similar for both types of materials and superconductivity should be achievable in n-type diamond provided a sufficiently large value of $N_e$ can be reached.

It is of interest to consider the possibility of heavily doped silicon becoming a superconductor through the same electron–phonon mechanism. Hole concentrations of the order of $N_h = 5 \times 10^{21}$ cm$^{-3}$ can be reached for B-doped Si. For $N_h = 3 \times 10^{21}$ cm$^{-3}$ Boeri et al. [3] have calculated $T_c = 0.3$ K, a temperature which depends strongly on the rather arbitrary value chosen for $\mu^*$ ($\mu^* = 0.1$). In order to repeat the simple calculation we have performed for diamond, we use a value of $d_0 = 40$ eV [24]. The corresponding density of states mass, however, is rather difficult to estimate with any degree of accuracy from the Luttinger parameters [28] using Eq. (1). A problem arises because the value of $m_\perp^{-1}$ \{1 1 0\} obtained in this manner for Si is close to zero within the uncertainties in the values of $\gamma_1$ [28,29]. Under these conditions $m_\perp$ \{1 1 0\} is close to infinity and the average performed by hand for diamond is not very meaningful for Si.

We presume that the same problem may afflict the results of ab initio calculations based on the local density approximation, because of the well-known gap problem. It would be of considerable interest to check for superconductivity p-type silicon samples as heavily doped as possible and at very low temperatures.

For n-type silicon doping of the order of $5 \times 10^{21}$ cm$^{-3}$ can also be reached. The density of state mass $\langle m^{1/2} \rangle = 1.1$ (2.5 for diamond), whereas the average electron–phonon coupling constant is about half of that for diamond [18,30]. The parameter $\lambda$ should thus be 0.06 times smaller than the corresponding one for diamond.
Since a rather small repulsion parameter $\mu^*$ would suffice to kill it, it is very unlikely that superconductivity will take place in n-type silicon.

The likelihood of superconductivity will be even smaller for germanium where the deformation potentials should be about the same as those for silicon, the effective density of states effective mass $4\langle m^{3/2} \rangle = 0.45$ i.e., less than half that of silicon.

Before closing this section, we mention that the possibility of superconductivity in n-type germanium was investigated early by Cohen [31]. For $N_c = 10^{20}$ cm$^{-3}$ (larger than the largest $N_c = 5 \times 10^{19}$ cm$^{-3}$ usually grown) he found a $T_c$ of a few millidegrees. He suggested that this $T_c$ would be larger ($T_c$ between 0.1 and 0.01 K) for a Ge–Si alloy with about 15% Ge because of the larger valley degeneracy (4 for Ge + 6 for Si for a Ge–Si alloy with about 15% Ge because of the possibility of superconductivity in n-type germanium.

The zero-point renormalization, by a single $T_c$ for the purpose of extrapolating to 0.25 K has, however, been observed in n-type (oxygen deficient) SrTiO$_3$ [32], a many-valley semiconductor with an energy gap of $\sim$3 eV, close to the lowest direct gap of silicon.

4. Electron–phonon renormalization of the absorption edge of semiconductors

It is well known that the electron–phonon interaction renormalizes the electronic states, in particular for semiconductors [5,13]. The oldest manifestation of this effect is a decrease in the absorption edge and the edge photoemission with increasing temperature. Fig. 3 shows the dependence of the free exciton cathodoluminescence of diamond on isotopic mass ($M = 12$ and 13). The peak energy is up-shifted by 13.6 meV when the $^{12}$C (which is almost exclusively (99%) found in natural diamond) is fully replaced by $^{13}$C (see caption of Fig. 3 for the contribution of the phonon frequency to this shift). This shift is related to the zero-point vibrational renormalization, which is proportional to $M^{-1/2}$ (see caption of Fig. 3). From the measured shift of 13.6 meV a total renormalization $\Delta E_g$ [33]:

$$\Delta E_g = 2 \times 13.6 \times M/\Delta M = 340 \text{ meV}$$

is found using the $M^{-1/2}$ dependence, $\Delta M = 1$ and $M = 12.5$. From the temperature dependence of $E_g$ a value $\Delta E_g = 370$ meV, in rather good agreement with Eq. (7) is estimated [5]. This rather strong renormalization should be of concern to computational theorists involved in ab initio calculations of the energy bands of semiconductors: these calculations assume that the atoms occupy fixed positions and thus neglect zero-point vibrations and the corresponding renormalizations.

For silicon ($\Delta E_g = 62$ meV) and germanium (52 meV) the zero-point renormalization is nearly an order of magnitude smaller (see Table 2 of [13]). This fact gave rise to the conjecture that electron–phonon interactions were particularly strong for semiconductors containing atoms of the second row of the periodic table (Be,B,C,N,O). Excellent candidates to test this conjecture are BN and SiC, but appropriate isotopic samples have not been measured. Isotopic mass and temperature dependence measurements have been recently performed for GaN [34] and ZnO [35].

In order to illustrate the determination of gap renormalizations by the electron–phonon interaction based on the temperature dependence of $E_g$, we display in Fig. 4 measurements of the A-exciton frequency vs. $T$ performed for GaN [34,36]. The temperature dependence of gaps or exciton energies is often represented, for the purpose of extrapolating to $T = 0$ and thus getting the zero-point renormalization, by a single Einstein oscillator (this procedure was followed for diamond in order to obtain the value of 370 meV mentioned above).

In the case of GaN (and also for ZnO), the phonons with predominant gallium vibrations have rather low frequencies, while those in which the oxygen motion dominates cluster around a much higher frequency. It is therefore necessary to use two oscillator frequencies.
The deformation potential \( d_0 \) for zinoblelde-type GaN (and also for ZnO) would thus be expected to lie around 60 eV. These materials would thus be more likely candidates than silicon for the observation of superconductivity, provided they can be prepared with a sufficiently high \( N_h \), an enormous challenge also of interest concerning electro-optic applications.

The question remains as to the physical reason for the high electron–phonon interaction in diamond and the corresponding nitrides and oxides. Concerning \( d_0 \), and also the gap renormalizations, both experiments [18,33] and calculations [13,17,38] confirm the large values. A hand-waving reason may be related to the absence of p-electrons in the core of the atoms of the second row of the periodic table. Under these conditions, the 2p valence electrons should move nearly rigidly with the vibrating cores and thus see fully the core potential. This would result in a strong electron–phonon interaction, especially for valence and conduction states involving 2p electrons.

We conclude this section by mentioning the case of ordered crystalline [39] and highly disordered [40] n-type InN. Thin films of this rather topical material, with \( N_e \) as high as \( 2 \times 10^{21} \text{ cm}^{-3} \), show a rapid decrease in resistivity for temperatures close to the \( T_c \) of metallic In (\( T_c \approx 3.4 \text{ K} \)). Although some reports attribute the superconductivity to heavily n-type InN (doped through nitrogen deficiency) [41], most of them leave the door open to the possibility of some form of metallic indium particles causing the superconductivity. One may, of course, conjecture that this superconductivity is related to that discussed in this work, especially in view of the high-electron concentrations of the superconducting InN sample and the presence of N (Which may induce strong electron–phonon interaction). However, the electrons in InN lie in a non-degenerate valley at the \( \Gamma \)-point of the conduction band, with an effective mass, \( m_e \approx 0.05 \) [42]. The corresponding density of states should, therefore, be much lower than that of the cases previously discussed. The most likely explanation is therefore the presence of some form of metallic indium [40].

5. Conclusions

The appearance of superconductivity in heavily borondoped diamond was firmly established at the time of the Tsukuba Workshop (December 2005). Most theoretical and experimental evidence points to a phonon-induced BCS mechanism. However, isotope effect measurements on \( T_c \), the smoking gun, are still missing. In this paper, we dissect the components of the BCS mechanism and suggest...
other materials in which superconductivity may (or may not) be present. Optical experiments which reveal the large electron–phonon interaction present in diamond (and also in GaN and ZnO) are discussed.

References

[1] E.A. Ekimov, V.A. Sidorov, E.D. Bauer, N.N. Mel’nik, N.J. Curro, J.D. Thompson, S.M. Stishov, Nature 428 (2004) 542.
[2] Y. Takanono, N. Nagao, I. Sakaguchi, M. Tachiki, T. Hatano, K. Kobayashi, H. Umezawa, H. Kawarada, Appl. Phys. Lett. 85 (2004) 2851.
[3] L. Boeri, J. Kortus, O.K. Andersen, Phys. Rev. Lett. 93 (2004) 237002.
[4] K.-W. Lee, W. Picket, Phys. Rev. Lett. 93 (2004) 237003.
[5] M. Cardona, Solid State Commun. 133 (2005) 3.
[6] Y. Ma, J.S. Tse, T. Cui, D.D. Klug, L. Zhang, Y. Xie, Y. Niu, G. Zou, Phys. Rev. B 72 (2005) 014306.
[7] H. Kato, S. Yamasaki, H. Okushi, Phys. Status Solidi (a) 202 (2005) 2122–2128.
[8] G. Baskaran, cond.-mat. 0410296, 2004.
[9] J. Nakamura, T. Oguchi, N. Yamada, K. Kuroki, K. Okada, Y. Takanono, cond.-mat. 0410144 (2004).
[10] K. Thonke, Semicond. Sci. Technol. 18 (2003) S20.
[11] J. Kačmarčík, C. Marcenat, C. Cytermann, A. Ferreira da Silva, L. Ortega, F. Gustafsson, J. Marcus, T. Klein, E. Gheeraert, E. Bustarret, Phys. Status Solidi (a) 202 (2005) 2160.
[12] L. Wei, P.K. Kuo, R.L. Thomas, T.R. Anthony, W.F. Banholzer, Phys. Rev. Lett. 70 (1993) 3764.
[13] M. Cardona, M.L.W. Thewalt, Rev. Mod. Phys. 77 (2005) 1173; H.J. Lian, A. Yang, M.L.W. Thewalt, R. Lauck, M. Cardona, Phys. Rev. B 28 (2006) 232302.
[14] P. Pavone, S. Baroni, Solid State Commun. 90 (1994) 295.
[15] M. Cardona, T. Ruf, Solid State Commun. 117 (2001) 201.
[16] R. Vogelgesang, A.K. Ramdas, S. Rodriguez, M. Grimsditch, T.R. Anthony, Phys. Rev. B 54 (1996) 3989.
[17] M. Cardona, N.E. Christensen, Solid State Commun. 58 (1986) 421.
[18] C. Canali, C. Jacoboni, F. Nava, L. Reggiani, S.F. Kozlov, in: Institute of Physics Conference Series No 43, 1979, p. 327 (Chapter 10).
[19] G. Abstreiter, M. Cardona, A. Pinczuk, Light Scattering in Solids IV: Electronic Scattering, Spin Effects, SERS and Morphic Effects, Springer, Berlin, 1984, p.5.
[20] J.W. Ager, III, W. Walukiewicz, M. McCluskey, M.A. Plano, M.I. Landstrass, Appl. Phys. Lett. 66 (1995) 616.
[21] M. Mermoux, A. Tajani, B. Marcus, E. Bustarret, E. Gheeraert, M. Nesladek, S. Koizumi, Diamond Relat. Mater. 13 (2004) 886.
[22] M. Willatzen, M. Cardona, N.E. Christensen, Phys. Rev. B 50 (1994) 18054.
[23] T. Yokoya, T. Nakamura, T. Matsuishi, T. Muro, Y. Takano, M. Nagao, T. Takenouchi, H. Kawarada, T. Oguchi, Nature 438 (2005) 647.
[24] A. Blacha, H. Presting, M. Cardona, Phys. Status Solidi (b) 126 (1984) 11.
[25] H.J. Xiang, Z. Li, J. Yang, J.G. Hou, Q. Zhu, Phys. Rev. B 70 (2004) 212504.
[26] X. Blase, Ch. Adessi, D. Connett, Phys. Rev. Lett. 93 (2004) 237004.
[27] E. Bustarret, J. Kačmarčík, C. Marcenat, E. Gheeraert, C. Cytermann, J. Marcus, T. Klein, Phys. Rev. Lett. 93 (2004) 237005.